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Ion acceleration with an ultra-intense two-frequency laser tweezer

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Abstract

Ultra-intense lasers produce and manipulate plasmas, allowing to locally generate extremely high static and electromagnetic fields. This study presents a concept of an ultra-intense optical tweezer, where two counter-propagating circularly polarized intense lasers of different frequencies collide on a nano-foil. Interfering inside the foil, lasers produce a beat wave, which traps and moves plasma electrons as a thin sheet with an optically controlled velocity. The electron displacement creates a plasma micro-capacitor with an extremely strong electrostatic field, that efficiently generates narrow-energy-spread ion beams from the multi-species targets, e.g. protons from the hydrocarbon foils. The proposed ion accelerator concept is explored theoretically and demonstrated numerically with the multi-dimensional particle-in-cell simulations.

1. Introduction

Laser acceleration of ions promises the compact source of energetic particles of interest for the basic science, industrial and medical applications [1, 2]. Today, with the modern high power lasers and advanced targetry, the record protons energies observed in the experiments approach 100 MeV [3–6]. In most of existing schemes, the accelerating fields are generated at the sheath of the laser-heated solid targets by the charge-separation between thermal electrons and the ion core [7]. In this case, the resulting ion flux is typically divergent and has a complex, broad spectrum of different ion species. In order to reach higher ion energies and beam quality, alternative acceleration mechanisms are actively studied, and to date include acceleration via radiation pressure [8, 9], collisionless shocks [10, 11], breakout afterburner [12], magnetic vortices [13–15], etc.

One way to improve acceleration efficiency is to use targets sufficiently thin, such that laser can push them via radiation pressure, thus accelerating all particles at the same rate, and the resulting ion flux is quasi-monoenergetic. This is know as the light sail (LS) acceleration [8, 16–19]. Despite being promising in an ideal one-dimensional (1D) geometry, it has turned out that in experiments maintaining the narrow ion spectra remains very challenging [20–22]. Besides the fact that LS process is rather demanding in terms of the laser and target quality, it is susceptible to a strong plasma instability. This basic process was observed both numerically [16, 17, 23, 24] and experimentally [25], and it can quickly destroy the target before the acceleration process is finished. Recent studies showed that this instability mainly results from the coupling between electron oscillations and the ion plasma modes [26, 27].

In this article we present a new acceleration concept, where two ultra-intense laser beams are used to spatially separate the electron and ion species acting as a relativistically intense optical tweezer.
In this approach, plasma instability is fully mitigated, and high gradient and high quality ion acceleration is produced by the optically controlled charge-separation. The proposed scenario relies on the two circularly polarized (CP) lasers of different frequencies colliding on a nano-foil, as shown in figure 1(a). Once lasers collide, their fields sum up into a beat wave, and its propagation direction and velocity are controlled by the ratio of the laser frequencies. Electrons of the foil are first squeezed into a thin sheet by the radiation pressure, and then are dragged by the beat wave, as if grabbed by a pair of tweezers (see figure 1(c)). This creates a nearly constant capacitor-like longitudinal electric field between the electron sheet and the slow ion core. For a case of the target containing a mixture of light and heavy ions, the lighter ions are quickly accelerated by this static field (see figure 1(b)).

In the following, we demonstrate that protons from a hydrocarbon foil can be accelerated by the tweezer to few hundreds of MeV on a distance of just a few micrometers. It is worth mentioning, that a similar principle is considered in the chirped-standing-wave acceleration [28], where electrons get trapped in the field of the laser reflected within a multi-layer nano-structure. Practically, such scenario requires both complicated targetry and highly chirped and precisely phased lasers, and electrons are not fully separated from the ions, so that LS instability may still persist. Compared to it, the proposed tweezer scheme operates with a single nano-foil, requires no specific phase control and provides a full separation of electrons and ions to avoid the plasma instability, making it simpler and more attractive for experimental realization.

2. Theory and numerical simulation

Demonstration presented in figure 1 is produced using a three-dimensional particle-in-cell (PIC) OSIRIS simulation [29], where two counter-propagating ten-cycle, temporally flat-top CP laser pulses with wavelengths $\lambda_1 = 0.8 \, \mu m$ and $\lambda_2 = 2\lambda_1$, and the same normalized amplitude $a_{1,2} = a_0 = 16$ are focused to $8 \, \mu m$ spots on a 20 nm hydrocarbon foil. The pre-ionized uniform plasma foil has electron density

![Figure 1](image-url)
n_e = 100n_c, proton density 10n_c, and C^{+6} ion density 15n_c, where n_c = 1.74 \times 10^{21} \text{ cm}^{-3} is the critical plasma density for 800 nm laser wavelength. Initial electron temperature is set to 1 keV. To resolve kinetics of the high density plasma, we use numerical grid with the longitudinal cell size \Delta z = 3 \text{ nm} and the transverse ones \Delta x = \Delta y = 10 \text{ nm}. To ensure low particle noise, 32 macro-particles per species per cell are used.

Let us now estimate the physical conditions required for the tweezer scenario. For the beat wave to be formed in the first place, the laser pulses (at least one) need to fully penetrate the thin foil. Based on the relativistic transparency, we can roughly define the laser strength requirement as \( a_0 \gtrsim 2\pi b_0 n_c / (\lambda_1 n_c) \), where \( b_0 \) is the initial target thickness. Note, that this condition is stronger than the standard relativistic transparency condition [30], as it ensures that plasma reflectivity is negligible compared to the transmission. This point is also validated with the PIC simulations and more details can be found in the supplementary materials.

Let us now consider electron dynamics in the electromagnetic wave created by two lasers with the frequencies \( \omega_1 > \omega_2 \), wavenumbers \( k_{1,2} \), and the same normalized amplitudes \( a_0 \). Presenting the lasers as two plane waves, we can write the resulting normalized vector potential as:

\[
a_{tw} = 2a_0 \sin k \cdot \xi_{slow}(e_y \cos k \cdot \xi_{fast} - e_x \sin k \cdot \xi_{fast}),
\]

where \( k = (k_1 \pm k_2)/2 \), laser wavenumbers \( k_{1,2} = \omega_{1,2}/c \), \( \xi_{slow} = (z - v_{tw} t) \) and \( \xi_{fast} = (z - c^2 t/v_{tw}) \). The value \( v_{tw} = (\omega_1 - \omega_2)/(k_1 + k_2) \) is a \'slow\' tweezer velocity. From the above equation, one can see that the beat wave field has two phase components \( k_1 \cdot \xi_{slow} \) and \( k \cdot \xi_{fast} \), which travel at the subluminal velocity \( v_{tw} \) and the superluminal velocity \( c^2/v_{tw} \), respectively.

Description of the relativistic motion of electrons in the beat wave equation (1) takes a simpler form in the reference frame moving with \( v_{tw} \) along the laser with shorter-wavelength. In this frame, both lasers have the same frequencies \( \omega_{tw} = \sqrt{\omega_1 \omega_2} \) and wavenumbers \( k_{tw} = \sqrt{k_1 k_2} \), and the total on-axis normalized vector potential can be written as,

\[
a_{tw} = 2a_0 \sin \left( k_{tw} z + \frac{\phi_0}{2} \right) \times \nabla \times \left[ e_y \cos \left( \omega_{tw} t + \frac{\phi_0}{2} \right) + e_x \sin \left( \omega_{tw} t + \frac{\phi_0}{2} \right) \right],
\]

where \( e_x \) and \( e_y \) are the transverse unit vectors, and \( \phi_0 \) is the carrier envelope phase difference of the two lasers. Then the longitudinal motion of a single electron follows the equation,

\[
\frac{dp_z}{dt} = -\frac{2a_0^2}{\gamma_e} m_e c \omega_{tw} \sin(2k_{tw} z + \phi_0),
\]

where \( p_z \) is electron longitudinal momentum in the co-moving frame, and the term in the right-hand side is the \( f_\perp \times B_\perp \) force of the CP standing wave. This force is focusing at the nodes and defocusing at the antinodes of the field, and electrons can get trapped around the nodes positions. Considering \( p_{z0} \) and \( z_0 \) to be initial longitudinal momentum and position of an electron, we can write its orbit in the \((z, p_z)\) phase plane as

\[
p_z^2 + 4a_0^2 m_e^2 c^2 \sin^2 [k_{tw}(z + z_0) + \phi_0] \sin[k_{tw}(z - z_0)] = p_{z0}^2.
\]

It is easy to see, that this trajectory is closed for \( |p_{z0}| \lesssim 2a_0 m_e c \), which introduces the electron trapping condition. In the plane wave, the transverse canonical momentum of electron is conserved, therefore the electron transverse momentum depends only on the phase of the standing wave (equation (2)), which for \( t = 0 \) gives \( p_{z0} = 2a_0 m_e c \sin (k_{tw} z_0 + \phi_0 / 2) \). Assuming that for a very thin foil \( z_0 = 0 \) for all particles, we can find their initial Lorentz factor \( \gamma_{z0} = \sqrt{1 + (p_{z0}/m_e c^2)^2} \approx 4a_0^2 \sin^2(\phi_0/2) / 2 \). By replacing \( p_{z0} \) with the initial longitudinal velocity \( v_{z0} \), the trapping condition, \( |p_{z0}| \lesssim 2a_0 m_e c \), can be rewritten as,

\[
\frac{v_{z0}^2}{c^2} \lesssim \frac{4a_0^2}{1 + 4a_0^2 \left[ 1 + \sin^2(\phi_0/2) \right]}. \tag{5}
\]

From equation (5) follows that for the ultra-intense lasers, \( a_0 \gg 1 \), electrons with initial velocities \( |v_{z0}| < c/\sqrt{2} \) can be trapped for any phase difference of the lasers. For the case of the interaction presented in figure 1, the initial velocities are mainly defined by the beat-wave velocity as, \( v_{z0} = -v_{tw} = -c/3 \), which is well within the trappable limit.
In order to account for the collective effects in the trapping process we should consider the charge-separation field (shown with arrows in figure 1(b)), which counteracts the \( v_\perp \times B_\perp \) force.

One can estimate the corresponding averaged Coulomb force as \( f_\text{c} = eE_{\text{sm}}/2 = 2\pi e^2 n_e l_0 \), where \( E_{\text{sm}} = 4\pi e n_e l_0 \) is the maximum charge-separation field after electrons are completely separated from the target. It is easy to see, that this force is actually small compared to the maximum of \( v_\perp \times B_\perp \) force, \( f_s \approx f_\text{c}/4 \), and the condition equation (5) is sufficient for the electron trapping. This point has been verified by the numerical simulations presented in supplementary material.

Once trapped, electron sheet moves helically in the laser field and travels with the beat-wave in a phase determined by both, laser \( v_\perp \times B_\perp \) force and static field of the ion plasma. Their averaged phase orbit is shown in figure 2(a). Since the transverse oscillations are relativistically strong, the electrons produce a broadband synchrotron light [31], which in this case is emitted into a ring-like shape at the angle \( \theta_{\text{emit}} = \arctan(\sqrt{c^2/v_\text{tw}^2 - 1}) \) to the tweezer propagation direction. In figure 2(b) we show the numerically reconstructed spectral-angular radiation distribution, where the central emission angle \( \theta = 66^\circ \) is close to theoretical estimate \( \theta_{\text{emit}} \approx 70^\circ \). The critical photon energy \( \hbar \omega_c \approx 0.25 \text{ keV} \), and total radiated energy reaches a few micro-Joules level for the trapped charge \( \approx 20 \text{ nC} \). See supplementary material for more details on the synchrotron light calculations.

For the considered interaction parameters, the charge separation field between electron sheet and ion plasma reaches \( E_z \approx 60 \text{ TV m}^{-1} \), and accelerates protons on a time scale, when the heavier carbon ions are essentially immobile (see figure 1(b)). This acceleration is nearly uniform, and it continues until protons outrun the electron sheet. The time at which a uniformly accelerated proton overruns the electron sheet \( t_{\text{max}} \) can be estimated as

\[
\int_0^{t_{\text{max}}} \frac{\alpha t}{\sqrt{1 + \alpha^2 t^2/c^2}} \, dt = v_{\text{tw}} t_m \Rightarrow t_m = \frac{2v_{\text{tw}} \gamma_{\text{tw}}^2}{\alpha},
\]

where \( \alpha = eE_z/m_p \) is the acceleration rate of protons, \( m_p \) is the proton rest mass, and \( \gamma_{\text{tw}} = c/\sqrt{c^2 - v_{\text{tw}}^2} \) is the Lorentz factor of the moving frame. Integrating the accelerated motion one can get the maximum energy gain of protons,

\[
\epsilon_{\text{max}} = 2m_p v_{\text{tw}}^2 \gamma_{\text{tw}}^2.
\]

In the simulation in figure 1, the acceleration process lasts for about 40 fs, and protons gain maximum energy of 220 MeV. These values are in good agreement with the theoretical estimate of 36 fs and 230 MeV given by equations (6) and (7) respectively. In figure 3(a), we show the proton phase space at the moment when they outrun the electrons. One can see a quasi-monoenergetic group of protons with an FWHM energy spread around 20% formed within a small divergence angle \( \theta_\perp \equiv P_\perp/P_\parallel \leq 10 \text{ mrad} \) (orange colormap and curve). This group contains \( \sim 10^{10} \) particles. Meanwhile, about \( 10^{11} \) protons with the broadband spectrum and larger divergence, \( \theta_\perp > 10 \text{ mrad} \), (green in figure 3) originate from the borders of irradiated region, where the accelerating field becomes three-dimensional (position dependent).

### 3. Discussion

For simplicity, the presented discussion and simulation assume the flat-top laser temporal profiles, which are hard to achieve in the experimental conditions. For this, we have run a series of additional 2D PIC simulations with the Gaussian temporal envelopes of the same pulse durations. This modeling
demonstrates, that the tweezer effect in this case is very similar to the idealized interaction, and production of 200 MeV protons in this case requires only slightly higher field amplitudes, $a_0 = 19$.

The finite transverse profiles of the laser beams also have an important effect on the scheme. Since the extracted electron sheet has a finite size, lasers can diffract at its edges, which erodes them and progressively reduces the size of the electron sheet. The time of sheet destruction via such mechanism $t_{de}$ should scale with the laser spot sizes and be of the order of $w_0/c$. From the simulations scanning the different laser radii, we could roughly estimate the scaling as $t_{de} \approx 1.3 w_0/c$. More details can be found in the supplementary material.

Based on these considerations, we can now estimate how the maximum proton energy scales with the input laser energy. Erosion time of the electrons from the sheet edge defines the required shortest laser durations as,

$$
\tau_{1,2} \simeq (1 \mp v_{tw}/c) t_{de},
$$

and their energies scale as,

$$
W_{1,2} \propto \tau_{1,2}(a_0 w_0)^2.
$$

Therefore, it gives the dependency of required laser energies on input wavelengths of two lasers, $W_1/W_2 = \lambda_2/\lambda_1$. Assuming that the optimal acceleration is reached when $t_{de}$ is close to the 1D maximum acceleration time $t_{max}$, one can obtain its relation from equation (6) as

$$
2v_{tw}^2 \gamma_{tw}^2 / \alpha = 1.3 w_0/c,
$$

and also the maximum proton energy as $\epsilon_{max} = 2m_p v_{tw}^2/\gamma_{tw}$ (from equation (7)). Besides, we choose $a_0 = 2\pi b_0 n_i/(\lambda_1 n_z)$ as the required minimal laser amplitude for relativistic transparency, which gives the longitudinal acceleration rate

$$
\alpha = eE_{em}/m_p = 4 \pi e^2 n_i b_0 / m_p = 2 a_0 \lambda_1 e^2 n_i / m_p.
$$

Based on all above mentions between $(W_{1,2}, \lambda_{1,2}, a_0, w_0$ and $\epsilon_{max})$, one can get the relation between the maximum proton energy with input total laser energy,

$$
\epsilon_{max}[\text{MeV}] \overset{\sim}{=} 0.34 \left( \frac{W_{\text{laser}} \lambda_1}{w_0} \right)^2 + 49.5 \frac{W_{\text{laser}} \lambda_1}{w_0},
$$

where the total laser energy $W_{\text{laser}} = W_1 + W_2$ is in Joules, and $\lambda_1$ and $w_0$ are in the same units. We note that the value of $\lambda_2$ is determined for different maximum proton energy that one wants to reach (see equation (7)).

We have tracked the scaling equation (8) in a series of 2D PIC simulations, by considering only the collimated protons, in $\theta_\perp \leq 10$ mrad (see figure 3(b)). In these simulations, both lasers have Gaussian temporal envelopes and super-Gaussian transverse profiles [e.g. $a \propto a_0 \exp(-r^2/w_0^4)$], and target ion composition is the same as in the above 3D case. The values of $\lambda_1 = 0.8 \mu m$ and $w_0 = 4 \mu m$ and $b_0 = 50$ nm are fixed, while other parameters were chosen in a way described previously in order to maintain the optimal acceleration conditions. As shown in figures 3(b) and (c), protons with 4%–20% energy spreads and energies from 20 MeV to 800 MeV are obtained for 1–80 J of laser energies. For the optimal conditions, $\lambda_2$ is chosen from 1 $\mu m$ to 3 $\mu m$ according to equation (7). This agrees well with the dependency given by equation (8). Assuming a cylindrical symmetry we can estimate the corresponding quantities of the physical protons in these simulations, which scale up from $0.3 \times 10^{10}$ to $2 \times 10^{10}$ for increasing laser energies.
For a more extensive comparison, in figure 3(c) we have plotted the 3D simulation and the results of the several 2D simulations of the ‘light sail’ scheme. For the latter points, we have considered a single laser, with the same parameters as the shorter wavelength tweezer laser, but higher laser intensity to match the same total energy. The target parameters are varying according to the optimal acceleration parameters given by [18, 19]. The result shows, that the present scheme can provide 2–5 times higher energy gain than the ‘light sail’ acceleration.

The robustness of the tweezer scheme against such experimental conditions as laser pulses time jitter and amplitudes imbalance was also verified numerically. These additional simulations are presented in the supplementary material (stacks.iop.org/NJP/22/052002/mmedia), and they demonstrate, that the acceptable time jitter is sentive to the differences of two lasers’ durations, and can be as high as 10 femtoseconds for current laser systems, and the normalized laser amplitudes should not deviate from each other more than, $|a_2/a_1 − 1| \leq 0.2$.

4. Conclusion

In conclusion, we have proposed a new scheme of ion acceleration from a nano-foil using two colliding relativistically intense lasers of different frequencies. The moving beat pattern of the lasers acts as a tweezer, that drags plasma electrons out of the ionized foil. The resultant capacitor-like field readily accelerates protons from the hydrocarbon plasma. The process is also accompanied by a strong synchrotron-like emission in the XUV/x-ray region, which can be used as a direct diagnostics of the interaction parameters, e.g. the tweezer phase velocity. In the experimental conditions, the two color laser pair, can be potentially produced using second-harmonic generation (SHG) [32] or OPCPA techniques [33, 34]. Thanks to recent techniques that improve temporal contrast of the laser, such as XPW [35], SHG [32], plasma mirrors [36] etc, this concept is now feasible experimentally. We anticipate this concept to become a promising candidate for the next generation of the compact laser-based high quality proton accelerators.

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