Inverse Compton scattering from solid targets irradiated by ultra-short laser pulses in the $10^{22} - 10^{23}$ W/cm$^2$ regime

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Received 17 December 2019, revised 13 March 2020
Accepted for publication 26 March 2020
Published 12 May 2020

Abstract
Emission of high energy gamma rays via the non-linear inverse Compton scattering process (ICS) in interactions of ultra-intense laser pulses with thin solid foils is studied using particle-in-cell simulations. It is shown that the angular distribution of the ICS photons has a forward-oriented two-directional structure centred at an angle $\vartheta = \pm 30^\circ$, a value which corresponds to a model based on a standing wave approximation to the electromagnetic field in front of the target, which only increases at the highest intensities due to faster hole boring, which renders the approximation invalid. The conversion efficiency is shown to exhibit a super-linear increase with the driving pulse intensity. In comparison to emission via electron-nucleus bremsstrahlung, it is shown that the higher absorption, further enhanced by faster hole boring, in the targets with lower atomic number strongly favours the ICS process.

Keywords: laser plasma, inverse Compton scattering, gamma rays, radiation reaction, foil targets, particle-in-cell

(Some figures may appear in colour only in the online journal)

1. Introduction

Next generation high-power laser systems are expected to routinely reach intensities in the $I \approx 10^{22} - 10^{23}$ W cm$^{-2}$ region [1–4]. In a configuration where such an intense pulse interacts with a solid target, gamma rays will be generated mostly by the processes of electron-nucleus bremsstrahlung [5], and by radiation reaction effects including non-linear inverse Compton scattering (ICS) [6, 7], where the fast electrons scatter on the high field of the laser pulse itself [8]. In this paper, we present a study of the latter process, relevant especially at the higher end of the considered intensity range, where the radiation has to be treated in the context of quantum electrodynamics (QED). The further outlook of even higher intensities would exhibit additional important effects such as the creation of electron-positron pairs and QED cascades [9–21], though we are not concerned with the pair-creation process throughout this paper which focuses solely on the emission of gamma rays in the interaction of laser photons with hot electrons. Apart from the aforementioned papers, which consider both ICS and pair-creation, several advanced schemes for all-optical gamma ray sources driven by a single laser pulse interacting with a solid target have been proposed [22–24]. In this paper, we describe the emission in a simplified situation of a short laser pulse interacting with a thin flat foil.

The non-linear multi-photon nature of the ICS process requires the presence of fast electrons and high fields. In the context of laser-plasma interactions, it has been observed in various configurations where the laser pulse interacts with an accelerated electron beam. Early observations [25–30] of the scattering of multiple photons on fast electrons were limited to the non-linear Thomson scattering regime characterized by low energy of the emitted photons, $h\omega_c \ll m_e c^2$, where $h$ is the reduced Planck constant, $\omega_c$ the photon’s angular frequency, $m_e$ the electron mass, and $c$ the speed of light [7]. Higher intensity lasers allowed for the observations
of the ICS interaction in the non-quantum regime, \( h\omega_\gamma \gg m_e c^2 \), in experiments with laser wakefield accelerated electrons and a counter-propagating laser pulse. These experiments generated gamma rays with energies \( \mathcal{E}_\gamma = 6 - 18 \) MeV [31], and \( \mathcal{E}_\gamma > 20 \) MeV [32], though the authors prefer to call the interaction the non-linear Thomson process in order to highlight that the quantum effects are still negligible in this regime. The energies sufficient to examine the quantum nature of the interaction were not reached until 2018 when a landmark experiment by Cole et al. [33] was performed at the Astra Gemini laser. It presented evidence of radiation reaction in the collision of an ultra-relativistic \( \mathcal{E}_e > 500 \) MeV electron beam, generated by laser-wakefield acceleration, with an intensity \( I = 1.3 \times 10^{21} \) \( \text{W cm}^{-2} \), \( \omega_0 = 25 \) laser pulse. The energy loss in the post-collision electron spectrum was correlated with the detected \( \mathcal{E}_\gamma > 30 \) MeV gamma ray signal, and was found to be consistent with a quantum description of radiation reaction. A further experiment [34] with a \( I = 4 \times 10^{20} \) \( \text{W cm}^{-2} \) pulse provided additional signatures of quantum effects in the electron dynamics in the external laser field, potentially showing departures from the constant cross field approximation in this intensity regime.

In contrast to the experiments where an intense laser pulse interacts with a solitary electron beam, the hot electrons participating in ICS in the laser-solid interactions studied in this paper are self-generated at the front side of the target due to the absorption [35, 36] of a portion of the energy of the same pulse with which they immediately interact, emitting high-energy gamma rays. By means of Particle-in-Cell simulations using the code EPOCH [37], we study the ICS emission from thin foils as a function of the laser pulse intensity, describe its energy spectrum and angular distribution, and present a simplified standing wave model that explains some of the emission’s prominent features. Additionally, we examine the effect of target material, and compare the ICS emission to bremsstrahlung, which we studied in our previous paper [38] under the same conditions.

The paper is organized as follows. Section 2 summarizes the essential theoretical background, and section 3 describes the PIC simulation setup. Section 4 presents the results, in particular the simulated ICS photon energy spectra, the simplified standing wave model and its comparison to the PIC simulations, the description of electron dynamics at the front side of the target, the predicted emission angle of the ICS photons and the angular distribution obtained from the PIC simulations, the efficiency of conversion of the driving laser pulse energy into that of the ICS photons, and a comparison of ICS to bremsstrahlung emission. Section 5 summarizes our conclusions.

2. Gamma ray emission by inverse Compton scattering

The ICS radiation is in fact not emitted continuously. Individual photons are emitted as the electron loses energy due to its interaction with the strong field. To characterize this interaction, taking into account the discontinuous nature of the process, a parameter \( \chi_e \) is introduced [6, 11, 39]:

\[
\chi_e = \frac{1}{E_5} \sqrt{\left(\gamma E + \frac{p \cdot B}{m_e}\right)^2 - \left(\frac{p \cdot E}{m_e}\right)^2}
\]

where \( E \) is the electric field, \( B \) is the magnetic field, \( p \) is the electron momentum, \( \gamma = 1/\sqrt{1-v^2/c^2} \) is the relativistic Lorentz factor of the electron, and \( E_5 \) is the “Sauter-Schwinger” field [40, 41], a critical field with enough strength to be able to perform \( m_e c^2 \) work over the electron Compton length \( \lambda_C = h/m_e c [11] \), \( E_5 = m_e c^3/e h = 1.32 \times 10^{16} \text{V cm}^{-1} \). Regarding the emission of gamma rays, the value of \( \chi_e \) indicates the strength of the radiation process, roughly separating the classical regime \( \chi_e \ll 1 \) with continuous emission, and the quantum regime, where \( \chi_e \) approaches unity and the process must be treated as a discontinuous emission of photon quanta [13, 42].

The intensity of the gamma radiation emitted by the electron can be expressed in the limits of \( \chi_e \ll 1 \) or \( \chi_e \gg 1 \) respectively as

\[
I_\text{rad}^{\leq c} = \frac{e^2 m_e^2}{6\pi e^2} \chi_e^2 (1 - c_1 \chi_e + c_2 \chi_e^2 - \ldots), \tag{2a}
\]

\[
I_\text{rad}^{> c} = c_3 \frac{e^2 m_e^2}{6\pi e^2} \chi_e^{2/3} (1 - c_4 \chi_e^{-2/3} + c_5 \chi_e^{-4/3} - \ldots) \tag{2b}
\]

where \( e \) is the elementary charge, and \( c_1, \ldots, c_5 \) are constants [43]. We can then give a rough estimate of the extreme limits for radiation intensity. At very small \( \chi_e \), we can only keep the unit term in the brackets of equation (2a), and the radiation intensity behaves as \( I_\text{rad} \sim \chi_e^2 \), while at very large \( \chi_e \), those terms in the brackets of equation (2b) which are inversely proportional to \( \chi_e \) raised to some positive power can be neglected, and the radiation intensity then behaves as \( I_\text{rad} \sim \chi_e^{-2/3} \).

Previous equations show that in order to generate large amounts of high energy gamma rays, one needs to employ a high field, hot electrons, or both. The strength of the laser pulse can be expressed in terms of the normalized amplitude of the vector potential

\[
a_0 = \frac{eE_0}{m_e\omega_c} \approx \left(7.3 \times 10^{-19} (\lambda [\mu m])^2 [W \text{cm}^{-2}]\right)^{1/2}, \tag{3}
\]

where \( E_0 \) is the peak amplitude of the electric field of the laser pulse, \( \omega_c \) its angular frequency, and \( \lambda \) its wavelength. The temperature of the hot electrons pulled out of a solid target by a pulse in the non-linear relativistic regime is given by

\[
T_e = m_e c^2 (\gamma - 1), \tag{4}
\]

the relativistic \( \gamma \) factor can be, in laser-solid interactions, approximated from the ponderomotive scaling [44] in the case of linear polarization as

\[
\gamma = \sqrt{1 + \frac{a_0^2}{2}}. \tag{5}
\]

For high values of \( a_0 \), this leads to a linear dependence \( T_e \sim a_0 \).
3. Simulation setup

Simulations were done in 2D in a \( x \in (-15, 15) \mu m \), and \( y \in (-20, 20) \mu m \) box with a cell size of 10 \( \times 10 \) nm. A normally incident laser pulse polarized in the simulation plane with a wavelength \( \lambda = 1 \mu m \), and a Gaussian spatial and temporal profile with a FWHM duration of \( \tau = 30fs \), was propagating along the \( x \) axis, and focused to a \( w = 3 \mu m \) spot at the front side of the target placed at \( x = 0 \). The laser pulse was emitted from the \( y = -15 \mu m \) boundary at the start of the simulation \( t = 0 \), at an angle of \( \theta_i = 0^\circ \) with its peak intensity crossing the boundary at \( t = 60fs \). The target was composed of a fully ionized CH plasma with electron density \( n_e = 289n_c \), where \( n_c = e\omega_0\omega/\epsilon^2 \) is the plasma critical density which is a function of the angular frequency \( \omega \) of the laser pulse, with \( \epsilon_0 \) being the permittivity of free space. For a \( \lambda = 1 \mu m \) laser pulse, the value of \( n_c = 1.1 \times 10^{21} cm^{-3} \). Parameter scans were performed for six laser pulse intensities between \( I = 3 \times 10^{21} \) W cm\(^{-2} \), and \( J = 10^{23} \) W cm\(^{-2} \). The normalized potential corresponding to the intensities in the simulations ranges from \( a_0 = 47 \) to \( a_0 = 270 \). Two additional materials, Al, and Au, were examined in order to compare these results to our previous work [38], which also describes their respective simulated parameters.

The simulations used a second order FDTD Maxwell solver [45], and a relativistic Boris pusher [46]. To limit noise and numerical heating [37], the simulations included a current smoothing algorithm and third order particle weighting. All boundary conditions were absorbing for radiation and thermalizing for particles. The radiation reaction effects were calculated with EPOCH’s Monte Carlo algorithm [47], and bremsstrahlung [38] was taken into account in order to obtain a self-consistent results. Electron-positron pair production processes were not taken into account, since they are only secondary, and would be exponentially suppressed in this intensity range. This paper only uses the photons with \( E_\gamma > 1MeV \) in all subsequent analysis.

The number of macro-particles varied along the \( y \) axis to ensure adequate resolution with \( N_{\gamma} = n_e/n_c \) electron macro-particles per cell in the middle of the simulated target, and save computational time at its far end where the background plasma dynamics is less violent. This was achieved by dividing the target into regions, schematically depicted in figure 1, with reduced number of particles per cell \( N_{\gamma} \) compared to the base value of \( N_{\gamma} = n_e/n_c \). To maintain the same initial electron density \( n_e \), these particles have been given an appropriately higher computational weight. Regions with lower \( N_{\gamma} \), summarized in table 1, were guarded with a thin layer of cells containing the base number of particles \( N_{\gamma} \) so that the simulated plasma expansion into the vacuum could represent densities lower than those represented by the higher weight particles. The target subdivision is the same as in [38] apart from the transverse extent of the target which is only \( y \in (-20, 20) \mu m \) in this paper. Unless explicitly stated otherwise, we model the target as an idealized flat surface foil with no presence of pre-plasma.

| Table 1. Number of macro-particles in different regions of the simulated target. The first column lists the designation \( N_{\gamma} \), of a computational region, which is equal to the number of electron macro-particles per cell in that region. The second column lists the value of \( N_{\gamma} \), i.e. the number of macro-particles per cell. The third column lists the number of ion macro-particles per cell either as an absolute value, or with respect to the number of electron macro-particles. The fourth column gives the extent of the given region along the \( y \) axis, and the last column shows the width of the ‘guard region’, composed of cells with the full number of macro-particles, at both ends of the \( x \) axis of the main region, where applicable. |
|---|---|---|---|
| region | number of electrons | number of ions | transverse extent (\( \mu m \)) | guard width (\( \mu m \)) |
| \( N_{\gamma} \) | \( n_e/n_c \) | \( n_e/Z \) | (0, 5) | – |
| \( N_{\gamma}/2 \) | \( n_e/2 \) | \( n_e/Z \) | (5, 10) | 0.3 |
| \( N_{\gamma}/10 \) | \( n_e/10 \) | \( n_e/10Z \) | (10, 15) | 0.2 |
| \( N_{10} \) | 10 | 2 | (15, 20) | 0.1 |

4. Results

4.1. Photon spectra

The EPOCH algorithm is set up so that the minimum energy of a photon emitted via the inverse Compton scattering mechanism is \( E_\gamma > 100keV \), though we limit the analysis to photons with \( E_\gamma > 1MeV \). In this case, there exists a threshold laser pulse intensity \( I \approx 3 \times 10^{21} \) W cm\(^{-2} \), below which no ICS-produced photons are seen in the simulation. At the lower intensities, as discussed in [38], the only photons with energies \( E_\gamma > 1MeV \) visible in the simulation are generated by bremsstrahlung. Photons generated by the bremsstrahlung mechanism are not included in subsequent analysis which focuses solely on the ICS mechanism, except for section 4.7 which summarizes the main differences between the radiation patterns resulting from the two mechanisms.

The tail of the distribution of the ICS-produced photons can be approximated by an exponential temperature fit \( N_{\gamma} \approx \exp(-E_\gamma/k_B T_\gamma) \), see figure 2. The effective temperature of
The electron temperature is dependent on the intensity as shown in figure 4. Though since the electric field strength varies with each laser cycle would ultimately lead to a Maxwell-Boltzmann like distribution. Thus, linear scaling of $\chi_e \sim a_0^2 \sim I$ follows straightforwardly from the theory, but the linear scaling of $T_\gamma \sim a_0^2 \sim I$, which turns out quite clearly in figure 3, should be treated as an empirical observation.

4.2. Standing wave model

The inverse Compton scattering process involves an electron moving in the field of the laser pulse in front of the target. To obtain more insight into the physical mechanisms governing the emission, we will make use of the simulation data with high temporal resolution with the help of a simplified theoretical model derived to describe the electron motion based on that of the electrons in our situation. As the immediate value of $\chi_e$ depends on the exact trajectory of the electron, a simple estimation for the most common energy of the resulting radiation has been proposed in the monochromatic approximation, giving $h\omega_\gamma \approx 0.44 \gamma m_e c^2$ [10, 13, 39]. This expression though describes the maximum of the photon distribution while the effective photon temperature $T_\gamma$ comes form a fit of the tail of a distribution which covers photons emitted by all of the electrons over the course of the simulation, therefore this expression cannot predict the temperature of the photons based on that of the electrons in our situation. As the immediate value of $\chi_e$ depends on the exact trajectory of the electron, a simple estimation connection between the temperature $T_e$ of the accelerated electron bunches and the temperature of the resulting radiation $T_\gamma$ cannot be made in the complex case of the laser-solid interaction where the bunch is of a finite size and, consequently, the different electrons interact with the field in a different phase. This is evident from the snapshot in figure 4, obtained from detailed studies of electron trajectories presented later in section 4.3, which shows the relation between the $\gamma$ factor and the $\chi_e$ parameter of the simulated electrons. We observe that there are many hot electrons which have the same $\gamma$ factor but span a broad range of attained $\chi_e$. Therefore, the immediate electron temperature $T_e$ does not readily reveal the radiation temperature $T_\gamma$, though averaging over many samples during the course of the whole interaction where both the $\gamma$ factor and the field strength vary with each laser cycle would ultimately lead to a Maxwell-Boltzmann like distribution.
on the standing wave approximation, which will be solved numerically.

As the electromagnetic wave of the linearly polarized laser pulse impinges on the highly overdense flat plasma slab at \( x = 0 \), most of it is reflected back and interferes with the incoming part of the pulse forming a standing wave in front of the target. The more equal the incident and reflected pulses, the more pronounced the standing wave pattern. In the case of a very short pulse, where the field intensity of the envelope changes rapidly with each oscillation, this pattern would be most prominent around the peak of the laser-target interaction where the intensity profile of the incoming and the reflected waves are close to being similar. This approximation ceases to be valid at extremely high intensities, where significant hole-boring and increased absorption leads to substantial destruction of the reflected wave-front [48, 49]. Though for the intensities studied in this paper, the deformity of the front side of the target is still low enough for the standing wave pattern to be sufficiently stable over the gamma ray emission timescale which is less than one half-cycle of the driving pulse.

The electric and magnetic field of the standing wave formed in front of the target in the case of normal incidence can be approximated by a plane wave near the interaction centre, and characterized by:

\[
\begin{align*}
E_y &= E_0 \sin(\omega t) \sin(kx), \\
B_z &= B_0 \cos(\omega t) \cos(kx),
\end{align*}
\]

where \( B_0 = E_0/c \). At the target’s surface, the \( E_y \) field then has a node, while the \( B_z \) field then has an anti-node. The maximum amplitude of the standing wave field is twice as large as that of the incident pulse due to the constructive interference of its incoming and outgoing parts.

In order to characterize the inverse Compton scattering radiation of an electron injected from the plasma surface into the standing wave, we expand equation (1) assuming \( \mathbf{B} = (0, 0, B_z) \) and \( \mathbf{E} = (0, E_y, 0) \). For high energy electrons with momenta \( p \gg m_e c \), we can make the approximation \( \gamma^2 \approx (p/m_e c)^2 \). Furthermore, as there are no forces acting on the electron along the \( z \) axis, \( p_z = 0 \), we can take \( \gamma^2 \approx (p_x/m_e c)^2 + (p_y/m_e c)^2 \). The last assumption allows to reorganize the terms in the expanded equation, leading to a simplified approximation

\[
\chi_e \approx \frac{1}{E_S} \left| \frac{p_x E_y}{m_e c} - c \gamma B_z \right|. 
\]

Equations (8) and (9) can be solved numerically, coupled with the relativistic equation of motion of the electron. Figure 5 shows the predicted maximum \( \chi_e \) attained by electrons of different initial momenta injected into the standing wave at different phase which radiate in the space in front of the target in the positive \( x \) direction. Around \( p_{x,0} = -20m_e c \) at a phase below \( \omega t_0 = \pi/2 \), there is a region of stability with respect to these two parameters in the sense that a small change of the initial conditions for the electron leads to a small change of maximum attained \( \chi_e \). Electrons injected with a much lower initial momentum do not radiate at all, while those with a much higher one will never return into the target, and will radiate in the backward direction. Such a high momentum injection cannot be achieved by the interaction of the laser pulse with the front side electrons, and does not appear in the full PIC simulations. However, similar trajectories, depicted in figure 6, can occur when recirculating electrons return form the back side of the target, and enter the area in front of the target while the pulse has a different phase than it would have had in case of direct injection from the front side. This kind of backward emission can be seen in very thin \( d \leq 2 \mu m \) foil in the late time of the interaction, being caused by the electrons which were injected early, and had enough time to do a subsequent full revolution in the target. Since the electron bunch spreads out in the transverse direction during the recirculation process [38], the returning electrons can be seen as essentially sampling arbitrary pulse phases in the \( (\omega t_0, p_x) \) phase-space.

For a sample numerical solution, shown in figure 7(a), we calculated the time evolution of the model for the initial momentum of \( p_{x,0} = -18m_e c \), which corresponds to the
Figure 6. Trajectories of electrons which were accelerated from the front side of a $d = 2 \, \mu m$ target during one of the early half-cycles of a $I = 10^{22} \, W \, cm^{-2}$ laser pulse. The simulation area depicted in the illustration spans approximately $x \in (-1, 0.5) \, \mu m$ and $y \in (-1, 1) \, \mu m$. The colour indicates the instant $\chi_e$ of the specific electron at that point with the scale going from blue (low) to magenta (high). Electrons were selected on the basis of attaining $\chi_e > 0.01$ during one laser pulse half-cycle, then their trajectories were plotted from the beginning of the half-cycle till the end of the simulation. The laser pulse was incoming from the left and injected many electrons into the target on an almost half-circle trajectory, seen in the lower right part of the picture. Upon entering the target, the electrons are not influenced by any strong fields, and continue in a straight line. After they reflect at the back side of the target (far right outside this illustration), some of them, albeit a much lower number, re-enter the interaction area with a high initial velocity, and radiate in the backward direction. This secondary emission happens at a late time of the interaction, and only those electrons that have had been injected in the earliest time arrive soon enough to meet the laser pulse at sufficient intensity to emit any significant amount of radiation – cf. total forward vs. backward emission in figure 13(a).

energy $\epsilon_e = 9 \, MeV$, injected into the $\pi/3$ phase of a standing wave which corresponds to the constructive interference of the incoming and reflected parts of an $a_0 = 86, I = 10^{22} \, W \, cm^{-2}$ laser pulse. The electron’s trajectory starts and ends at the surface of a target positioned at $x = 0$. The model tracks the evolution of the electric $E_y$ and magnetic $B_z$ fields along the trajectory of the simulated electron. Together with the electron’s $p_x$ and its $\gamma$ factor, these constitute the two parts of the simplified equation (9). The calculation shows that as the electron moves through the evolving field of the standing wave, it is first accelerated in the $y$ direction, giving rise to its $\gamma$ factor with little change to its negative $p_x$. After the phase-reversal of $B_z$, the $p_x$ starts to rise until the electron turns around and starts to be accelerated in the $+x$ direction around the moment where $E_y$ reaches its minimum. The highest emission parameter $\chi_e$ is then attained right before the re-injection into the target, where the negative $B_z$ reaches its minimum, and $E_y$ is almost negligible.

The result in figure 7(a) compares favourably to the actual trajectory of an electron selected form the PIC simulation on the basis of similar injection phase, and the initial and final relativistic $\gamma$ factor, in figure 7(b), showing that the standing wave model is a good approximation to the trajectory of a single electron.
4.3. Electron dynamics

In order to describe the dynamics of the electrons responsible for the gamma ray emission via inverse Compton scattering, we can compare the results of the numerical solution of equations (8) and (9), seen in figure 7(a), to a simulation snapshot zoomed-in to the centre of the interaction area in figure 8. It shows the trajectories of a random sample of electrons which achieve a high value of $\chi_e$ during one half-cycle of the driving laser pulse. In the simulation, a total of about 18 000 electron macro-particle reach $\chi_e > 0.01$ during this particular half-cycle, and over 99% of them follow trajectories of a similar shape as the one produced by the aforementioned sample numerical solution. At this stage of the interaction, hole boring by the laser pulse has pushed the target surface from $x = 0$ to $x \approx 150\text{nm}$, the phase of the $E_y$ field is changing, and a new bunch is about to be accelerated.

First, an electron is pulled out of the target surface, and injected into the standing wave in front of the target when the balance between the $J \times B$ force and the force due to the $E_x$ field is violated. This stage is not covered by the theoretical model, where instead inject an electron with a specified initial momentum (or a range of momenta, as will be described in the following text), and neglect the $E_x$ field altogether.

After being injected, the electron is accelerated in the $+y$ direction by the $E_y$ field, causing a rise in its relativistic $\gamma$ factor. Meanwhile, the phase of the $B_z$ field changes, causing the increase in the originally negative momentum $p_x$ up to a moment when $p_x = 0$, and the electron is at the maximum distance $\Delta x \approx 180\text{nm}$ away from the actual target surface.

Next, the rising $B_z$ field transforms the transverse momentum $p_y$ into the longitudinal $p_z$ as the electric field $E_y$ weakens. The relativistic $\gamma$ factor is dominated by the $p_z$ component — in the normalized units of figure 7, $p_z \approx \gamma$, and the electron is returning into the target with $p_x \gg p_y$.

Maximum $\chi_e$ parameter is attained right before the re-injection, when $\gamma$ is almost constant as $E_y$ is decreasing with the impending phase change. The $c\gamma B_z$ is now the dominant term in equation (9), but due to the still non-negligible $p_y$, maximum emission occurs at an angle $\alpha \neq 0$. Right before the re-injection, the $B_z$ field starts to decrease, and the electrons, which have lost most of their transverse momentum continue to propagate inside the target. The process is about to repeat with the forthcoming laser pulse half-cycle, albeit mirrored with respect to the $x$ axis.

4.4. Emission angle

As we have seen that the maximum emission occurs when the electron is propagating at an angle, we shall now discuss some features of the angular distribution of the emitted photons seen in the theoretical model. Figure 9 shows that the theoretical model predicts an angle $\alpha_{\text{max}}$, measured from the $x$ axis, where the emission parameter $\chi_e$ has a maximum for an electron with a given initial momentum. To see how the
Figure 9. Time evolution of the emission parameter $\chi_e$, and the propagation angle $\alpha$, measured for the $x$ axis, of the emitting electron in the simplified theoretical model with the initial momentum corresponding to $\varepsilon_e = 9\text{MeV}$ injected into the $\omega t_0 = \pi/3$ phase of a standing wave formed by the reflection of a $I = 10^{22} \text{W cm}^{-2}$ pulse.

Figure 10. Theoretical model of the maximum emission parameter $\chi_e$, and the corresponding emission angle $\alpha$ for different initial energies of electrons injected into the $\omega t_0 = \pi/3$ phase of a standing wave formed by the reflection of a $I = 10^{22} \text{W cm}^{-2}$ pulse.

Figure 11. Theoretical model of the maximum emission parameter $\chi_e$ reached by any electron for a given laser pulse potential $a_0 = \sqrt{I}$ (where the standing wave maximum intensity is $I_{SW} = 2I$) with the angle $\alpha_{max}$ at which the emission occurs, the relativistic factor $\gamma_{max}$ attained by the electron at the point of maximum emission, and the relativistic factor $\gamma_0$ with which has the electron been injected into the standing wave.

Figure 12. The most prominent direction of propagation of the high energy photons described as the mode of the angular distribution of all photons with energies above the 50th percentile for $d = 2 \mu\text{m}$ foils from CH and aluminium for driving pulses of different intensities. For $a_0 = 50$, the 50th percentile corresponds to $E_\gamma \sim 1\text{MeV}$, while for $a_0 = 270$, the limit is $E_\gamma \sim 3\text{MeV}$.

The angle of maximum emission changes in case when a spectrum of electrons would be injected, we first calculate the model values for a range of initial electron momenta. For each energy, we find the time $t_{max}$ when the emission parameter has a maximum $\chi_{e,\text{max}} = \chi_e(t_{max})$, $d\chi_e/dt|_{t_{max}} = 0$, and the angle $\alpha_{\text{max}} = \alpha(t_{max})$ at which the maximum emission occurs for the given electron energy. Figure 10 shows that there is an optimal initial electron energy $\varepsilon_{e,\text{opt}}$ which leads to the highest value of the emission parameter at a given laser pulse intensity. Electrons around this optimum are responsible for the majority of the gamma radiation, while those which are too far away, be they slower or faster, would emit considerably less.

Then, we perform a parameter scan over laser pulse intensities, finding the optimal initial electron energy $\varepsilon_{e,\text{opt}}$ which leads to the highest value of the emission parameter at a given laser pulse intensity. Electrons around this optimum are responsible for the majority of the gamma radiation, while those which are too far away, be they slower or faster, would emit considerably less.

Our assumptions hold, one can expect this to be the direction of maximum emission of the ICS gamma rays in the simulations.

4.5. Angular distribution

In the PIC simulations, the angular distribution of photons emitted via the inverse Compton scattering process in the interaction with a $I = 10^{22} \text{W cm}^{-2}$ pulse has a distinct structure with two lobes centred around $\vartheta \simeq 30^\circ$ and $\vartheta \simeq 330^\circ$. This result is consistent both with previously published simulations [8, 50, 51], and the theoretical model presented in section 4.3.

In the case of very thin foils $d < 2\tau_{L}$, recirculating electrons have enough time to make a full revolution and return to the front side of the target while the interaction with the laser pulse is still ongoing. This then leads to an appearance of backward radiation, which is suppressed for thicker foils. The $d = 2 \mu\text{m}$
Figure 13. Angular distribution of photons emitted via the inverse Compton scattering process from CH foils with \( d = 2 \, \mu \text{m} \) at different driving pulse intensities \( I = 10^{22} \, \text{W cm}^{-2} \) (a) and \( I = 10^{23} \, \text{W cm}^{-2} \) (b). The different curves represent the sum of the energies of all photons in respective energy span in the units of conversion efficiency of the total laser pulse energy into gamma rays in that energy span in given direction per \( 1^\circ \) shown on the radial axis. The selection of energy bands in the figures here is not fixed, but differs between simulations to represent exclusive percentile ranges, indicated in the figure legend, to highlight the similarities of the structure of the spectra which, for different intensities, appear at different absolute energy values.

The geometry of the front side is defined by the hole boring process \([52]\) since we do not observe any significant decoupling \([14]\) of the ion and electron fronts. As the plasma is being pushed forward, the depth of the ion front increases gradually in the transverse direction towards the centre forming an angled side-wing which stretches from near the focus centre at \( y = 0 \), where the hole reaches the maximum depth, to the region with much lower pulse intensity several micrometers away from the centre, where the original target surface is virtually undisturbed.

As the intensity increases, faster hole boring leads to a larger incidence angle at the sides of the hole, and we cannot assume that the electrons are pulled in front of the target in the direction normal to the polarization of a standing wave. Instead, some enter the interaction area at higher angles. While the radiation is still predominantly forward-going even for the highest intensity \( I = 10^{23} \, \text{W cm}^{-2} \) examined in this paper, with increasing intensity, the emission angle increases, backward radiation is enhanced, and the shape of the resulting spectrum, shown in figure 13(b), is approaching that of “transversely oscillating electron synchrotron emission” (TOEE) \([51]\), which itself, in simulations parametrised on plasma density, can be seen as an intermediate stage between the emission from a highly overdense \([13]\) and a near-critical-density \([53, 54]\) target. Detailed exploration of such low density regimes is out of scope of this paper, nevertheless the highest-intensity case presented here bears some similarity to the TOEE process. Furthermore, in this high-intensity short pulse interaction, carrier envelope phase effect leads to a pronounced asymmetry of the emitted radiation.

While the hole boring process influences the gamma ray angular distribution in the case of a solid foil with a flat surface, an even more profound effect is revealed in simulations which include pre-plasma, where the interaction moves to a regime of a laser pulse propagating through underdense plasma. This stage is characterized by side injection from a higher density plasma edge formed by electrons pushed away by the ponderomotive force into positively charged channel. Energy stored in the space charge field is then released as periodic pulses of backwards propagating electrons which are in turn slowed by the radiation reaction force \([53]\) and emit high energy photons in the backward direction. This process is called ‘reinjected electron synchrotron emission’, or RESE \([14]\). For an exponential pre-plasma profile with the scale length of \( l = 1 \, \mu \text{m} \), the trajectories of the electrons injected
from the lower density regions are chaotic, as seen in figure 14, with no readily identifiable typical features. When the laser pulse reaches the overdense target, hole boring and reflection occur as in the case without pre-plasma, emitting a similar spectrum with the angular distribution featuring the two forward lobes at approximately ±30°. The resulting angular distribution, shown in figure 15 is a combination of both processes. Moreover, since the electrons are accelerated to higher energies in lower density plasma, the emission is enhanced even in the forward direction, where it retains the original structure.

4.6. Conversion efficiency

Figure 16 shows that in our simulations, the total conversion efficiency obeys the scaling

\[ \eta_{ICS} \sim I^{1/2} \]  

for both the aluminium and the CH targets. Similar efficiency dependence has been observed in other simulations [50]. As we have established, in equation (7), the emission parameter scales linearly with the laser pulse intensity, \( \chi_e \sim I \). According to equations (2a) and (2b), the gamma radiation intensity scales as \( I_{rad} \sim \chi_e^{2} \) with the power \( \zeta = 2 \) for \( \chi_e \gg 1 \), and \( \zeta = 2/3 \) for \( \chi_e \ll 1 \). Our simulations reach up to \( \chi_e \approx 1 \), a region where neither of the proposed limits are valid. On the one hand, should we lower the intensity to attain \( \chi_e \ll 1 \), no ICS emission would be seen at all. On the other, with much higher intensities where \( \chi_e \gg 1 \) would be attained, we can no longer speak about an interaction with an opaque over-critical target because of the onset of relativistic transparency. Since we have \( \chi_e \sim I \), equation (10) suggests that the region in question could be reasonably described by an intermediate empirical value of \( \zeta = 3/2 \).

4.7. Comparison to Bremsstrahlung

In an experiment, the detectors themselves cannot distinguish between the gamma rays emitted due to bremsstrahlung, which we explored in a previous paper [38], and those emitted due to the inverse Compton scattering process studied here. Both will be seen at the same time, and the distinction has to be based on distilling their unique features from the total spectra.
The first question to be answered is whether the radiation generated by the respective processes would be seen at all. In figure 17(a), we see that for CH foils, ICS dominates already at the lowest intensity $I = 5 \times 10^{21}$ W cm$^{-2}$ where it is detectable. Both its temperature and the number of generated photons rise quickly with the rising intensity, much faster than that of bremsstrahlung. The combination of a thin low-Z target and the inverse Compton scattering process does not depend on the target thickness, therefore the same values can be used for comparison of the $d = 2 \mu m$ and the $d = 5 \mu m$ foils.

![Figure 17.](image.png)

**Table 2.** Efficiency of conversion of the laser pulse energy into the energy of all photons generated by the bremsstrahlung process $\eta_{BS}$, and the inverse Compton scattering process $\eta_{ICS}$ for targets of different materials irradiated by a $I_0 = 86, I = 10^{22}$ W cm$^{-2}$ laser pulse. For foils with $d > 2 \mu m$, the conversion by inverse Compton scattering does not depend on the target thickness, therefore the same values can be used for comparison of the $d = 2 \mu m$ and the $d = 5 \mu m$ foil.

| material      | thickness | $\eta_{BS}$ | $\eta_{ICS}$ |
|---------------|-----------|-------------|--------------|
| C$^{6+}$H$^+$ | $2 \mu m$ | 1.6         | 690          |
|               | $5 \mu m$ | 3.5         |              |
| Al$^{13+}$    | $2 \mu m$ | 5.2         | 240          |
|               | $5 \mu m$ | 12          |              |
| Au$^{51+}$    | $2 \mu m$ | 89          | 75           |
|               | $5 \mu m$ | 190         |              |
| Au$^{10+}$    | $5 \mu m$ | 140         | 31           |

5. Conclusions

We have studied the emission of gamma rays by inverse Compton scattering in interactions of a short intense laser pulse with a thin foil target via 2D PIC simulations. The ICS process dominates over bremsstrahlung in low-Z targets already at a threshold intensity $I \approx 3 \times 10^{21}$ W cm$^{-2}$ under which no ICS generated gamma rays are seen at all. Spectra of the gamma rays produced in interactions with different driving pulse intensities show a linear dependence of the ICS produced gamma ray temperature on the intensity $T_{\gamma} \sim a_0^2 \sim I$, at least in the studied intensity range $I \approx 3 \times 10^{21} - 10^{23}$ W cm$^{-2}$.

The radiation is forward going with two lobes centred at approximately $\vartheta \approx \pm 30^\circ$. The angular distribution of the emission is dictated by the dynamics of the electrons in the field of the laser pulse in front of the target, thus for sufficiently thick $d \gtrsim 2 \mu m$ targets, there is no change in its structure with increasing thickness.

A simple theoretical model which assumes the movement of an electron in a planar standing wave formed in the front side by the interaction of the incoming and reflected parts of the laser pulse predicts the photon propagation angle $\vartheta = 30^\circ$ regardless of the laser pulse intensity. This is confirmed by the simulations up to $I \approx 10^{22}$ W cm$^{-2}$. As the intensity grows further, the propagation angle increases since the assumptions of the theoretical model break down due to hole boring. When the hole in the surface is sufficiently deep, the electrons injected from its sides meet the laser pulse in a different phase, and travel along a different trajectory before being reinjected near the centre of the hole.

Efficiency of conversion of the driving laser pulse energy into that of the gamma rays generated by ICS shows super-linear scaling with intensity $\eta_{ICS} \sim I^{3/2}$ in the studied intensity range.

Comparing the results to our previous work, where we show that targets made of materials with a higher atomic number, while exhibiting a lower absorption, still show a significant increase of gamma ray production by bremsstrahlung [38], we see that the lower absorption also affects the ICS process which does not directly depend on the atomic number. The efficiencies of both processes are roughly equal in the case of a $d = 2 \mu m$ Au$^{51+}$ target irradiated by a $I = 10^{22}$ W cm$^{-2}$ laser pulse.
Acknowledgments

We would like to express our thanks to Mariana Kecová from the ELI Virtual Beamline team for the visualization of electron trajectories in figure 6. The results of this work were obtained under Project LQ1606 with the financial support of the Ministry of Education, Youth and Sports as part of targeted support from the National Programme of Sustainability II. Supported by the project ELITAS (CZ.02.1.01/0.0/0.0/16 013/0001793) and the project HiFi (CZ.02.1.01/0.0/0.0/15 003/0000449) from the European Regional Development Fund. Additional funding was obtained under GACR project 18-09560S, and the SGS16/248/OHK4/3T/14 grant. The EPOCH code used in this research was developed under UK Engineering and Physics Sciences Research Council grants EP/G054940/1, EP/G055165/1 and EP/G056803/1. Simulations were performed at the ECLIPSE cluster at ELI-Beamlines, supported from the aforementioned ELI related grants, and the Salomon cluster at IT4Innovations, supported by the Ministry of Education, Youth and Sports from the Large Infrastructures for Research, Experimental Development and Innovations project “IT4Innovations National Supercomputing Center – LM2015070”.

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References

[1] Weber S et al 2017 Matter Radiat. Extremes 2 149–76
[2] Hernandez-Gomez C et al 2010 J. Phys.: Conf. Ser. 244 032006
[3] Zou J P et al 2015 High Power Laser Sci. Eng. 3 2052–3289
[4] Danson C, Hillier D, Hopps N and Neely D 2015 High Power Laser Science and Eng. 3 2052–3289
[5] Bernstein M J and Comisar G G 1970 J. Appl. Phys. 41 729–33
[6] Ritus V I 1985 J. Sov. Laser Res. 6 497–617 1573-8760
[7] Lau Y Y, He F, Umstadter D P and Kowalczyk R 2003 Phys. Plasmas 10 2155
[8] Nakamura T, Koga J K, Esirkepov T Z, Kando M, Korn G and Bulanov S V 2012 Phys. Rev. Lett. 108 195001
[9] Gu Y J, Jirka M, Klimo O and Weber S 2019 Matter Radiat. Extremes 4 064003
[10] Bell A R and Kirk J G 2008 Phys. Rev. Lett. 101 1079–7114
[11] Bulanov S V, Schroeder C B, Esarey E and Leemans W P 2013 Phys. Rev. A 87 1094–1228
[12] Di Piazza A, Müller C, Hatsagortsyan K Z and Keitel C H 2012 Rev. Mod. Phys. 84 1177–228
[13] Ridgers C P, Brady C S, Duclos R, Kirk J G, Bennett K, Arber T D, Robinson A P L and Bell A R 2012 Phys. Rev. Lett. 108 1079–7114
[14] Brady C S, Ridgers C P, Arber T D, Bell A R and Kirk J G 2012 Phys. Rev. Lett. 108 245006
[15] Zhidkov A, Koga J, Sasaki A and Uesaka M 2002 Phys. Rev. Lett. 88 185002
[16] Luo W, Liu W Y, Yuan T, Chen M, Yu J Y, Li F Y, Sorbo D D, Ridgers C P and Sheng Z M 2018 Sci. Rep. 8 1–8
[17] Luo W, Wu S D, Liu W Y, Ma Y Y, Li F Y, Yuan T, Yu J Y, Chen M and Sheng Z M 2018 Plasma Phys. Control. Fusion 60 095006
[18] Jirka M, Klimo O, Bulanov S V, Esirkepov T Z, Gelfer E, Bulanov S S, Weber S and Korn G 2016 Phys. Rev. E 93 023207
[19] Grismayer T, Vranic M, Martins J L, Fonseca R A and Silva L O 2017 Phys. Rev. E 95 023210
[20] Yu J Y, Yuan T, Liu W Y, Chen M, Luo W, Weng S M and Sheng Z M 2018 Plasma Phys. Control. Fusion 60 044011
[21] Luo W et al 2015 Phys. Plasmas 22 063112
[22] Ta Phuoc K, Corde S, Thaury C, Malka V, Tafzi A, Goddet J P, Shah R C, Sebben S and Rousse A 2012 Nat. Photon. 6 308–311 1749-4893
[23] Gong Z, Hu R H, Lu H Y, Yu J Q, Wang D H, Fu E G, Chen C E, He X T and Yan X Q 2018 Plasma Phys. Control. Fusion 60 044004
[24] Zhu X L, Chen M, Yu T P, Weng S M, Hu L X, McKenna P and Sheng Z M 2018 Appl. Phys. Lett. 112 174102
[25] Englert T J and Rinehart E A 1983 Physical Review A 28 1539–45
[26] Bula C et al 1996 Phys. Rev. Lett. 76 3116–19
[27] Chen S Y, Maksmchuk A and Umstadter D 1998 Nature 396 653–5
[28] Schwoerer H, Liesfeld B, Schlenvoigt H P, Amthor K U and Sauerbrey R 2006 Phys. Rev. Lett. 96 1079–7114
[29] Malka G et al 2002 Physical Review E 66 066402
[30] Chen S et al 2013 Phys. Rev. Lett. 110 155003
[31] Sarri G et al 2014 Phys. Rev. Lett. 113 223401
[32] Yan W et al 2017 Nat. Photon. 11 514–20
[33] Cole F J et al 2018 Phys. Rev. X 8 011020
[34] Poder K et al 2018 Phys. Rev. X 8 031004
[35] Malka G and Miquel J L 1996 Phys. Rev. Lett. 77 75–8
[36] Pukhov A 2001 Phys. Rev. Lett. 86 3562–5
[37] Arber T D et al 2015 Plasma Phys. Control. Fusion 57 113001
[38] Vyskočil J, Klimo O and Weber S 2018 Plasma Phys. Control. Fusion 60 054013, 1561-6587
[39] Kirk J G, Bell A R and Arka I 2009 Plasma Phys. Control. Fusion 51 085008
[40] Sauter F 1931 Z. Phys. 69 742–64
[41] Schwinger J 1951 Phys. Rev. 82 664–79
[42] Shen C S and White D 1972 Phys. Rev. Lett. 28 455–9
[43] Nikishov A and Ritus V 1985 Sov. Phys. JETP 19 529–41
[44] Wilks S C, Kruer W L, Tabuk M and Landgon A B 1992 Phys. Rev. Lett. 69 1383–6
[45] Yee K 1996 IEEE Trans. Antennas Propag. 14 302–7
[46] Boris J P 1970 Relativistic plasma simulation—Optimization of a hybrid code Proc. 4th Conf. on Numerical Simulations of Plasmas (Naval Research Laboratory, Washington DC, 1970) pp 3–67
[47] Duclos R, Kirk J G and Bell A R 2010 Plasma Phys. Control. Fusion 53 015009
[48] Yuan T, Yu J Y, Liu W Y, Weng S M, Yuan X H, Luo W, Chen M, Sheng Z M and Zhang J 2018 Plasma Phys. Control. Fusion 60 065003
[49] Liu W Y, Luo W, Yuan T, Yu J Y, Chen M and Sheng Z M 2017 Phys. Plasmas 24 103130
[50] Ji L L, Pukhov A, Nerush E N, Kostyukov I Y, Shen B F and Akli K U 2014 Phys. Plasmas 21 023109
[51] Chang H X, Qiao B, Zhang Y X, Xu Z, Yao W P, Zhou C T and He X T 2017 Phys. Plasmas 24 043111
[52] Robinson A P L, Gibbon P, Zepf M, Kar S, Evans R G and Belleci C 2009 Plasma Phys. Control. Fusion 51 024004
[53] Brady C S, Ridgers C P, Arber T D and Bell A R 2013 Plasma Phys. Control. Fusion 55 124016
[54] Brady C S, Ridgers C P, Arber T D and Bell A R 2014 Phys. Plasmas 21 031018