Attosecond gamma-ray pulses via nonlinear Compton scattering
in the radiation dominated regime

Jian-Xing Li, Karen Z. Hatsagortsyan, Benjamin J. Galow, and Christoph H. Keitel
Max-Planck-Institut für Kernphysik, Saupfercheckweg 1, 69117 Heidelberg, Germany

The interaction of a relativistic electron bunch with a counter-propagating tightly-focused laser beam is investigated for intensities when the dynamics is strongly affected by its own radiation. The Compton scattering spectra of gamma-radiation are evaluated employing a semiclassical description for the laser-driven electron dynamics and a quantum electrodynamical description for the photon emissions. We show for laser facilities under construction that gamma-ray bursts of few hundred attoseconds and dozens of mega-electronvolt photon energies may be detected in the near-backwards direction of the initial electron motion. Tight focussing of the laser beam and radiation reaction are demonstrated to be jointly responsible for such short gamma-ray bursts which are independent of both duration of electron bunch and laser pulse. Furthermore, the stochastic nature of the gamma-photon emission features signatures in the resulting gamma-ray comb in the case of the application of a multi-cycle laser pulse.

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Shortly after the invention of the laser, it was realized that Compton scattering [1] of laser radiation by a relativistic electron beam can be a bright source of x- and gamma-rays [2–3]. Later proof-of-principle experiments [4–9] showed picosecond hard x-rays with electron beams from a linear accelerator. With appearance of the laser wakefield acceleration technique for electrons, the all-optical setup for Compton/Thomson radiation sources from a few hundreds of keV up to 8–9 MeV photon energies, with a shorter duration of about 50 fs laser pulse duration, has been elaborated theoretically [10–15] and experimentally [16–25]. While these experiments are based on linear Compton scattering, which allows for narrowband gamma-radiation sources for nuclear resonance fluorescence [26–27], recently efforts succeeded in demonstrating Thomson scattering in the nonlinear regime [28].

In superstrong laser fields Compton scattering acquires nonlinear characteristics due to multiple laser photon absorption [29–33]. Moreover, radiation reaction can enter into play for these extreme conditions [38–37]. The radiation-dominated regime (RDR), when the radiation reaction has a decisive impact on the electron dynamics, is characterized by the parameter $R \equiv \alpha \xi \chi \gtrsim 1$, [37], where $\alpha$ is the fine structure constant, $\xi \equiv |e|E_0/(m\omega_0)$ the invariant laser field parameter with $E_0$ and $\omega_0$ being the laser field and frequency, respectively, $\chi \equiv \gamma(\omega_0/m)(1 - \beta \cos \theta)$ the quantum strong field parameter with $\beta$ and $\gamma$ being the relativistic Lorentz factors of the electron, $\theta$ the polar angle of the electron velocity with respect to the laser propagation direction, and $e$ and $m$ are the electron charge and mass, respectively (Planck units $\hbar = c = 1$ are used throughout). The RDR is most accessible in the quantum regime of interaction when $\chi \gtrsim 1$, [37]. As RDR regime is attainable with present petawatt laser systems and taking into account the prospects connected with next generation lasers (see e.g., [38–40]), a lot of attention recently has been devoted to investigations of radiation reaction effects in Compton scattering [41–49].

In linear Compton scattering the duration of the resulting gamma-radiation pulse is determined by the shortest duration of the laser and electron beams and, usually, limited by the electron bunch length. The nonlinearity of the process and focusing of the laser pulse can further assist in shortening of the generated pulse [50]. Nevertheless, multi-MeV gamma-ray pulses of less than femtosecond duration are still not available, while they could be indispensable for time-resolved nuclear spectroscopy [51, 52].

In this Letter we investigate the feasibility of generating multi-MeV gamma-rays of several hundreds of attoseconds duration via nonlinear Compton scattering of a tightly focused laser pulse off a counterpropagating elec-

![FIG. 1. (Color online) Schematic scenario for the considered generation of ultrashort gamma-ray bursts which arise from a relativistic electron beam counterpropagating with a superstrong laser pulse. The front electrons of the electron beam loose sufficient energy due to radiation reaction to be reflected and to emit brief gamma-ray bursts when leaving the laser focal region.](https://arxiv.org/abs/1504.02393v2)
electron beam in the quantum RDR. Both, the tight focusing of the laser beam and the essential influence of the radiation reaction are necessary ingredients to induce a reflection of a fraction of the electron bunch and a burst of radiation within a half-cycle of the laser field, i.e., ultrashort gamma-photon pulses. Due to the essential nonlinear mechanism, the duration of the produced gamma-ray pulses (or duration of the pulse in the train) is independent of the laser and electron beam durations. We implement Monte-Carlo simulations of the electron radiation based on QED, while propagating the electrons between photon emissions according to classical equations of motion \[53–55\]. For the ultrashort gamma-ray production the following laser and electron parameters are required: \(R = \alpha \chi \gtrsim 1\) and \(\chi \approx 10^{-6} \gamma \xi \lesssim 1\) for realizing quantum RDR, and \(\gamma \sim \xi/2\) to allow for electron reflection, which finally requires \(\xi \sim \gamma/10^{3}\). The emitted gamma-photon energy is proportional to \(\chi\), subsequently to \(\xi\). Note that all employed parameters are within the achievable limit of current or next generation lasers, see e.g., \[33, 43\].

We employ a linearly polarized short focused laser pulse, which is an approximate solution of Maxwell equations with respect to the diffraction parameter \((k_0 w_0)^{-1}\) and the temporal parameter \((\omega_0 T_0)^{-1}\) \[43\], where \(k_0, w_0\) and \(T_0\) are the wave vector, the waist radius and the pulse duration of the laser beam, respectively.

In superstrong laser fields \(\xi \gg 1\), the coherence length of the photon emission is much smaller than the laser wavelength and the typical size of the electron trajectory \[33\]. Then, the photon emission probability \(W\) is determined by the local electron trajectory, consequently, by the local value of the parameter \(\chi\) \[33, 43\]:

\[
\frac{d^2 W}{d\eta d\omega} = \frac{\alpha \chi m^2}{\sqrt{3\pi}} \int_0^\infty \frac{K_{5/3}(x)dx}{\tilde{\omega}_r \tilde{\chi}^2 K_{2/3}(\omega_r)},
\]

where \(\eta = \omega t - k_0 z\), \(\tilde{\omega} = k_0^\mu \cdot k_\mu / (\tilde{\chi} k_0^\mu \cdot p_\mu)\) is the normalized emitted photon energy, \(\tilde{\chi} = 3\chi/2\), \(k_0^\mu\) and \(p^\mu\) are the four-vectors of the driving laser photon, the emitted photon and the electron, respectively, and \(\tilde{\omega}_r = \tilde{\omega} / p_0\) with recoil parameter \(p_0 = 1 - \tilde{\chi}\) (in the classical limit \(p_0 \approx 1\)). The photon emission of electrons is considered to be a Monte-Carlo stochastic process \[33\]. Between the photon emissions, the electron dynamics in the laser field is governed by classical equations of motion: \(dp/dt = -e(E + v \times B)\). Given the smallness of the emission angle \(\sim 1/\gamma\) for an ultrarelativistic electron, the photon emission is assumed to be along the electron velocity. The photon emission induces the electron momentum change \(p_f \approx (1 - \omega/cp)p_i\), where \(p_{i,f}\) are the electron momentum before and after the emission, respectively, and \(\omega\) is the emitted photon energy. During a small step of propagation \(\Delta \eta\), the photon emission will take place if the condition \(dW/d\eta|_{\Delta \eta} \geq N_e\) is fulfilled, where \(N_e\) is a uniformly distributed random number in \([0, 1]\) (the value of \(\Delta \eta\) is chosen small enough to keep the total number of photon emissions consistent). The photon energy \(\omega\) is determined by the relation: \(1/W \int_{\omega_{min}}^{\omega} (dW(\omega)/d\omega) d\omega = \tilde{N}_r\), where \(\tilde{N}_r\) is an another independent random number in \([0, 1]\), and, \(\omega_{min}\) is the minimal energy of the emitted photon, restricted by the laser photon energy. The radiation intensity is defined as the emission energy per unit detector time \(t_d = t - \mathbf{n} \cdot \mathbf{r}(t)/c\), where \(\mathbf{n}\) is the radiation direction and \(\mathbf{r}\) the electron coordinate.

We consider radiation when an ultrarelativistic electron bunch with an initial kinetic energy \(K_i = 200\) MeV (an initial energy \(\gamma_0 \approx 392\)), counterpropagates an ultrashort focused laser pulse in the quantum RDR. The schematic set-up is shown in Fig. 1. The laser peak intensity is assumed \(I \approx 4.9 \times 10^{23}\) W/cm\(^2\), the laser wavelength \(\lambda_0 = T_0 = 1\) \(\mu m\), and the laser beam waist size \(w_0 = 1\) \(\mu m\) (\(\xi = 600\), and \(\chi_{max} \approx 0.8\)). The electron reflection condition in the laser field should hold \(\gamma(t_r) \approx \xi/2\) with reflection time \(t_r\), which requires larger electron initial energies \(\gamma_0 > \xi/2\) because of radiation losses. The electron bunch length here is \(l_b = 10 \lambda_0\), and the transverse size \(w_b = w_0\). The energy as well as angular spread of the bunch is \(\Delta \gamma/\gamma_0 = \Delta \theta = 10^{-3}\), the number of electrons in the bunch \(N_e = 3 \times 10^{8}\), and the electron density \(n_b = 3 \times 10^{19}\) cm\(^{-3}\). The space-charge force \(F_{cc} \sim 2a_0 n_b\).
is negligible with respect to the laser force $F_L \sim \xi m \omega_0$ in this case, as $F_C/F_L \sim 2 \times 10^{-4}$, where $x_b$ is the electron beam extension along the laser field.

The gamma-radiation properties above the photon energy of 1 MeV in a 4-cycle laser pulse are shown in Fig. 2. The time-dependent angular resolved radiation intensity in Fig. 2(a) shows that the emission is mostly significant near the emission polar angles $\theta \approx 180^\circ$ and $\theta \approx 20^\circ$. The radiation at $\theta \approx 180^\circ$ is straightforward to understand following the dynamics of a single electron in the bunch as illustrated in Fig. 2. The ultrarelativistic electrons in the bunch radiate in the forward direction mostly near the strong field region $\chi \approx \gamma \xi \sim 1$ before they reach the reflection point (see Fig. 2(a)), which is indicated by the large change of $\eta$ at $\eta \approx 36$ in Fig. 3(d). The duration of the emission is determined solely by the length of the electron bunch. In fact, this can be estimated via $\tau_d \approx \tau - n_{150^\circ} \cdot r(t)/c \approx \tau + z_b/c$, where $\tau$ is the emission duration in the electron time with $\tau \approx z_b/c = 10T_0$ and laser period $T_0$ according to Fig. 2(b), which shows the radiation intensity vs the electron time. Consequently, this emission duration is governed in time by the length of the electron bunch via $\tau_d \approx 2z_b/c = 20T_0$.

The radiation near $\theta \approx 20^\circ$ is, however, governed by a much shorter time scale for two reasons. During the forward reflection all electrons lose energy which facilitates the electron reflection when the condition $\gamma \sim \xi/2$ is approached. For our chosen timings and parameters, then only the front fraction of the bunch can meet the strongest laser field in the maximally focussed region and loose sufficient energy to fulfill the reflection condition. Due to laser defocusing rear electrons merely experience weaker laser fields and consequently will not achieve the reflection condition. Furthermore considering those electrons in the reflected front fraction, significant photon emission occurs when the $\chi$-parameter reaches the rather large peak value $\chi \approx 0.6$ (see Fig. 2(a)) which decreases the electron energy (see Fig. 2(c)). Subsequently, the discussed reflection happens at $\eta \approx 36$. After the reflection such electrons emit backwards. This emission [at $\theta \approx 20^\circ$ around $\eta \approx 37$ in Fig. 3(d)] is even brief on the scale of a single cycle, because for the indicated angle those electrons rapidly escape the focal region of the laser beam (with $\{\xi, \chi\} \rightarrow 0$). This is confirmed by Fig. 2(b) displaying that the duration of the backward emission of the bunch in the electron time is $\tau \approx \tau_0 < t_b$. This duration is then drastically shortened in the detector time $\tau_d \approx \tau(1 - \beta \cos \theta)$, yielding the considered case at $\theta \approx 20^\circ$ an estimated $\tau_d \approx 0.25T_0$.

The time duration of the gamma-radiation at $\theta = 20^\circ$ with an aperture angle $\Delta \theta = 0.002$ is illustrated in Fig. 2(c). The duration of the main gamma-pulse is about $\tau_\gamma = 0.25T_0 = 830$ as. The corresponding spectral distribution is shown in Fig. 2(d), with the central frequency being $\omega \approx 67m = 34.2$ MeV. Narrower gamma-ray pulses can be detected at smaller polar angles $\theta$, where, however, the central frequency is smaller, since $\chi$ is smaller. Similar spectra can be detected also at the azimuthal angle $\phi \approx 0$.

Several gamma-ray bursts are observable near $\theta \approx 20^\circ$ in the 4-cycle laser pulse (see Fig. 2(a)). Moreover, a single gamma-ray burst arises in a 2-cycle laser pulse, and a gamma-ray comb is formed by employing longer laser pulses. The number of bursts in the comb is determined by the number of laser periods in the laser pulse (see Fig. 4), because the electron can be reflected at any wave crest. Due to the stochastic character of the gamma-photon emission, there is a probability that the electron will emit forwards a large amount of energy (and be reflected) not in the first crest of the laser field but in the next ones. Additionally, the different electrons in the bunch can be reflected in different laser cycles. The role of the stochasticity can be seen comparing Fig. 2(a) and (b). When stochastic effects are neglected, there is only...
one single burst observable near $\theta \approx 20^\circ$. This feature of the angle-resolved radiation intensity can serve as an indicator of stochastic effects in photon emission.

The ultrashort gamma-ray pulses in Fig. 2 are induced by the combined effect of laser focusing and radiation reaction as analyzed in Fig. 5. In the simulation of Fig. 5(a), the radiation dominated regime is considered, however, rather than a focused laser pulse as in Fig. 2 a plane-wave 4-cycle Gaussian pulse is employed. As there is no defocusing effect, all electrons in the bunch experience the same strong laser field and are reflected. Three bands can be seen in Fig. 5(a), which correspond to the radiation from the electrons reflected in adjacent laser cycles near the peak of the laser pulse. After a reflection the corresponding electron moves along the laser pulse propagation direction, experiencing the field of the rest of the laser pulse and emitting photons during the long time of the order of $\tau_0/(1-\beta)$, which yields a radiated pulse of the rather long duration of $\tau_0(1-\beta \cos \theta)/(1-\beta)$. Consequently, the backward emission is broad in the plane wave laser field. Figure 5(b) illustrates the case out of the radiation dominated regime with $\xi = 100$ and $\gamma_0 = 39.2$ ($\chi \approx 0.01$) and using a focused laser beam. In this case the radiation loss is quite small and the reflection condition is assumed to hold for the initial electron energy $\xi \sim 2\gamma_0$. Therefore the reflection can take place in any cycle of the laser pulse. The emission duration is then determined mostly by the total laser pulse duration.

Furthermore, we investigate the question which part of the electron bunch is most responsible for the generation of the short gamma-ray bursts. In Fig. 6 the angle-resolved radiation intensity is shown with respect to the initial coordinate of the electrons in the bunch. Figure 6(a) confirms that the front electrons are reflected due to radiation reaction ($z \lesssim 4\lambda_0$), but the rear electrons do not because of the laser defocusing effect. The backward radiation along $\phi = 0$ strongly declines when the electron’s initial $y$-coordinate in the bunch does not deviate much from the center of the beam, $y \lesssim 0.001\lambda_0$ (see Fig. 6(b)). This is mainly due to the $y$-component of the laser field which is nonvanishing in the focused linearly polarized laser beam, increasing linearly with $y$-coordinate variation from the center, and diverts the radiation into a large azimuthal angle. When carrying out the radiation simulations, we took advantage of the mentioned property, and used a thin slice of the electron bunch with the number of electrons calculated from the given density of the electron bunch.

Finally, we estimate the total number of photons in the gamma-ray burst of 830 as duration for the parameters in Fig. 2 to be $N_{\text{ph}} \sim 3 \times 10^8$ within the emission solid angle $\Delta \Omega = 1 \text{ mrad}^2$. In spite of a small number of total photons, the photon flux (F) and the brilliance (B) are rather large due to the short duration of the pulse: $F \sim 4 \times 10^{15}$ photons s$^{-1}$, 0.1%BW, and $B \sim 10^{22}$ photons s$^{-1}$ mrad$^{-2}$ mm$^{-2}$ 0.1%BW, respectively, i.e., the brilliance is 4 orders of magnitude larger than in the recent experiment 25.

Concluding, we have shown that brilliant attosecond gamma-ray bursts can be produced by the combined effect of laser focusing and radiation reaction in nonlinear Compton scattering in the radiation dominated regime. A gamma-ray comb is formed when applying a long laser pulse, which carries signatures of stochastic effects in photon emissions.

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