Signatures of the Unruh Effect from Electrons Accelerated by Ultrastrong Laser Fields

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We calculate the radiation resulting from the Unruh effect for strongly accelerated electrons and show that the photons are created in pairs whose polarizations are perfectly correlated. Apart from the photon statistics, this quantum radiation can further be discriminated from the classical (Larmor) radiation via the different spectral and angular distributions. The signatures of the Unruh effect become significant if the external electromagnetic field accelerating the electrons is not too far below the Schwinger limit and might be observable with future facilities. Finally, the corrections due to the birefringent nature of the QED vacuum at such ultrahigh fields are discussed.

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Introduction.—One of the most fascinating phenomena of noninertial quantum field theory is the Unruh effect: An observer or detector undergoing a uniform acceleration a experiences the Minkowski vacuum as a thermal bath with the Unruh temperature [1]

$$T_{\text{Unruh}} = \frac{\hbar}{2\pi k_B c} a. \tag{1}$$

As one might expect from the principle of equivalence, the Unruh effect is closely related to Hawking radiation [2], i.e., black hole evaporation: The uniformly accelerated observer (detecting the Unruh effect) corresponds to an observer at a fixed distance to the horizon (feeling the gravitational pull and measuring the Hawking radiation), whereas the inertial observer in flat space-time is analogous to an unfortunate astronaut freely falling into the black hole. However, there is also a crucial difference between the two phenomena: In contrast to the case of uniform acceleration, the free fall into a black hole is (per the definition of a black hole) not invariant under time reversal. Hence, while Hawking radiation generates a real outflow of energy (black hole evaporation), the Unruh effect corresponds to an equilibrium thermal bath and does not create any energy flux *per se*.

The most direct way of observing this striking effect would be to accelerate a detector and to measure its excitations. However, this is extremely difficult since moderate accelerations correspond to extremely low temperatures, and, thus, the Unruh effect has not been directly observed so far (see, however, [3,4]). Therefore, we focus on a somewhat indirect signature in the following: Since the uniformly accelerated detector acts as if it was immersed in a thermal bath, there is a finite probability that it absorbs a (virtual) particle from this bath and passes to an excited state. Translated back into the inertial frame, this process corresponds to the emission of a real particle [5]. The opposite process, when the detector reemits the particle into the bath in the accelerated frame and goes back to its ground state, also corresponds to the emission of a real particle in the inertial frame.

In the limiting case that the time between absorption and reemission becomes arbitrarily small, the detector transforms into a scatterer which scatters particles from one mode into another mode of the thermal bath in the accelerated frame. Translated back into the inertial frame, this process corresponds to the emission of two real particles by the accelerated scatterer. This effect is analogous to moving-mirror radiation [6] and can be interpreted as a signature of the Unruh effect. In the following, we calculate this quantum radiation given off by electrons accelerated in ultraintense electromagnetic fields (acting as pointlike noninertial scatterers) and compare it to the classical (Larmor) radiation. An analogous idea has already been pursued in Ref. [7] but in the derivation presented therein did not take into account crucial features of the radiation (such as the fact that the photons are always created in correlated pairs).

Low-energy effective action.—For accelerations caused by electric fields $\boldsymbol{E}$ well below the Schwinger [8,9] limit (the regime we are interested in), the Unruh temperature (1) and hence also the typical energies of the scattered photons are much smaller than the electron mass m. Furthermore, we assume that the magnetic field $\boldsymbol{B}$ is negligible compared to the (external) electric field $\boldsymbol{E}$ accelerating the electron: $\boldsymbol{E}^2 \gg c^2\boldsymbol{B}^2$. Apart from facilitating a straight electron trajectory (which will become important later on), this ensures that the energy $\boldsymbol{\mu}_e \cdot \boldsymbol{B}$ associated to the spin $\boldsymbol{\mu}_e \approx \boldsymbol{\sigma} e\hbar/m$ of the electron is much smaller than the photon energies. Hence, the spin only changes very slowly and basically does not interact with the other degrees of freedom. In this limit, the electrons can be treated as pointlike particles which act via their charge only (Thomson scattering). In the temporal gauge, the corresponding action $\mathcal{A} = -m\int ds + q\int dx^\mu A_\mu$ yields the

Lagrangian ($\hbar = \varepsilon_0 = \mu_0 = c = 1$ throughout)

$$L = -m\sqrt{1-\dot{\boldsymbol{r}}^2} - q\dot{\boldsymbol{r}}\cdot\boldsymbol{A}(\boldsymbol{r}), \tag{2}$$

where q is the charge of the electron with the trajectory $\boldsymbol{r}(t)$ and $\boldsymbol{A}$ the vector potential.

Since it will be most relevant for the following investigations, we shall focus on planar Thomson scattering (which yields the maximum cross section), where the polarizations $\boldsymbol{A}_{\rm in}$ and $\boldsymbol{A}_{\rm out}$ of the scattered photons are perpendicular to the (undisturbed) trajectory of the electron. Hence, the field $\boldsymbol{A}$ in (2) can be split up $\boldsymbol{A} = \boldsymbol{A}_\parallel + \boldsymbol{A}_\perp$ into a strong external electric field $\boldsymbol{E}_\parallel = \partial_t \boldsymbol{A}_\parallel \approx$ const accelerating the electron plus weak electromagnetic waves $\boldsymbol{A}_\perp$ with low energies. They cause small deviations $\delta\boldsymbol{r}_\perp$ of the electron's path $\boldsymbol{r} = \boldsymbol{r}_\parallel + \delta\boldsymbol{r}_\perp$ from its linear trajectory $\boldsymbol{r}_\parallel$. After linearization, the transversal canonical momentum $\delta\dot{\boldsymbol{r}}_\perp/(1-\dot{\boldsymbol{r}}_\parallel^2)^{1/2} - q\boldsymbol{A}_\perp/m$ is conserved along the electron's path according to the Euler-Lagrange equations, and, thus, we may eliminate the small transversal fluctuations $\delta\boldsymbol{r}_\perp$ by inserting $\delta\dot{\boldsymbol{r}}_\perp \approx q(1-\dot{\boldsymbol{r}}_\parallel^2)^{1/2}\boldsymbol{A}_\perp/m$ back into the action (2). Consequently, the dynamics of the weak electromagnetic waves $\boldsymbol{A}_\perp$ is governed by the low-energy effective action for planar Thomson scattering

$$\mathcal{L}_\perp = \tfrac{1}{2}(\boldsymbol{E}_\perp^2 - \boldsymbol{B}_\perp^2) - g\boldsymbol{A}_\perp^2\,\delta^3(\boldsymbol{r}_\parallel[t]-\boldsymbol{r})\sqrt{1-\dot{\boldsymbol{r}}_\parallel^2[t]}. \tag{3}$$

In the following, we shall drop the superscripts $\perp$ for brevity and denote the (undisturbed) electron's trajectory by $\boldsymbol{r}_e = \boldsymbol{r}_\parallel$. The coupling $g = q^2/m$ determines the cross section for planar Thomson scattering, and the last factor ensures the relativistic invariance of the action $\mathcal{A} = g\int ds A_\mu A^\mu - \int d^4x F_{\mu\nu}F^{\mu\nu}/4$. In the natural units used here, the charge q is related to the fine-structure constant $\alpha_{\rm QED}$ via $q = \sqrt{4\pi\alpha_{\rm QED}} \approx 0.3$.

The effective action (3) reproduces the well-known picture of Thomson scattering: For weak electromagnetic waves whose wavelength is much larger than the Compton wavelength (formal limit $m \uparrow \infty$), the electron acts as a classical pointlike scatterer with a spin and energy-independent cross section [9]. Since the electron is too heavy to feel the recoil of the scattered photons, the correlations between the photons and the path of the electron (and also the time delay between absorption and reemission mentioned in the introduction) are negligible. Finally, for planar Thomson scattering, the angular dependence of the scattering amplitude vanishes.

Quantum radiation.—In order to calculate the photons created by the noninertial motion $\boldsymbol{r}_e[t]$ of the scatterer, let us split the total Hamiltonian into a perturbation part

$$\hat{H}_1(t) = g\hat{\boldsymbol{A}}^2(t,\boldsymbol{r}_e[t])\sqrt{1-\dot{\boldsymbol{r}}_e^2[t]}, \tag{4}$$

supplemented with the usual adiabatic switching on and off $g(|t|\uparrow\infty) = 0$, and the undisturbed Hamiltonian $\hat{H}_0 = \frac{1}{2}\int d^3r(\hat{\boldsymbol{E}}^2 + \hat{\boldsymbol{B}}^2)$, which leads to the usual normal mode expansion.

Since the coupling $g \approx 3.5\times 10^{-14}$ m is much smaller than all other relevant length scales (such as the wavelengths of the photons), higher orders in g can be neglected and the evolution of the initial Minkowski vacuum $|0\rangle$ can be derived via time-dependent perturbation theory $|{\rm out}\rangle = |0\rangle - i\int dt\hat{H}_1(t)|0\rangle + \mathcal{O}(g^2)$, which gives

$$|{\rm out}\rangle = |0\rangle + \sum_{\boldsymbol{k},\lambda,\boldsymbol{k}',\lambda'}\mathfrak{A}_{\boldsymbol{k},\lambda,\boldsymbol{k}',\lambda'}|\boldsymbol{k},\lambda,\boldsymbol{k}',\lambda'\rangle + \mathcal{O}(g^2), \tag{5}$$

with the two-photon amplitude

$$\mathfrak{A}_{\boldsymbol{k},\lambda,\boldsymbol{k}',\lambda'} = \frac{\boldsymbol{e}_{\boldsymbol{k},\lambda}\cdot\boldsymbol{e}_{\boldsymbol{k}',\lambda'}}{2iV\sqrt{kk'}}\int dt g\sqrt{1-\dot{\boldsymbol{r}}_e^2[t]} \times \exp\{i(k+k')t - i(\boldsymbol{k}+\boldsymbol{k}')\cdot\boldsymbol{r}_e[t]\}. \tag{6}$$

Here $\boldsymbol{k}$ is the wave number and $k = |\boldsymbol{k}|$ the frequency of the photon modes; λ labels their polarization described by the unit vector $\boldsymbol{e}_{\boldsymbol{k},\lambda}$, and V is the quantization volume. Since a time-resolved detection of the created photons is probably infeasible, polarization and momentum are the best observables to be measured. As one may infer from the above expression, the photons are always emitted in pairs (squeezed state), and there is a perfect correlation of the polarizations of the two photons due to the scalar product $\boldsymbol{e}_{\boldsymbol{k},\lambda}\cdot\boldsymbol{e}_{\boldsymbol{k}',\lambda'}$. Parallel photons $\boldsymbol{k}\parallel\boldsymbol{k}'$ are maximally entangled; i.e., they must have the same polarization $\lambda = \lambda'$. Note that this applies to linear polarization; the circular polarizations of the two created photons are opposite (for Thomson scattering [9]) due to angular momentum conservation.

In terms of the new integration variable $\tau = t - r_e^\parallel[t]$ with $r_e^\parallel = (\boldsymbol{k}+\boldsymbol{k}')\cdot\boldsymbol{r}_e/(k+k')$, the two-photon amplitude

$$\mathfrak{A}_{\boldsymbol{k},\lambda,\boldsymbol{k}',\lambda'} = \frac{\boldsymbol{e}_{\boldsymbol{k},\lambda}\cdot\boldsymbol{e}_{\boldsymbol{k}',\lambda'}}{2iV\sqrt{kk'}}\int d\tau g\frac{\sqrt{1-\dot{\boldsymbol{r}}_e^2[t]}}{1-\dot{r}_e^\parallel[t]}e^{i(k+k')\tau}, \tag{7}$$

is determined by the Fourier transform of the effective (direction-dependent) Doppler factor above.

Larmor radiation.—In order to discuss the observability of this quantum radiation, it must be compared with the competing classical process. The Larmor radiation can be derived from the relativistic action $q\int dx^\mu A_\mu$ and corresponds to a coherent state $\propto \exp\{\sum_{\boldsymbol{k}}\alpha_{\boldsymbol{k},\lambda}\hat{a}^\dagger_{\boldsymbol{k},\lambda}\}|0\rangle$ with the coefficients (see, e.g., [10])

$$\alpha_{\boldsymbol{k},\lambda} = q\int dt\frac{\boldsymbol{e}_{\boldsymbol{k},\lambda}\cdot\dot{\boldsymbol{r}}_e[t]}{\sqrt{2Vk}}e^{ikt - i\boldsymbol{k}\cdot\boldsymbol{r}_e[t]}. \tag{8}$$

The numerator $\boldsymbol{e}_{\boldsymbol{k},\lambda}\cdot\dot{\boldsymbol{r}}_e$ displays the well-known blind spot in the forward and backward direction $\boldsymbol{k}\parallel\dot{\boldsymbol{r}}_e$. The introduction of a new integration variable $\tau = t - \boldsymbol{k}\cdot\boldsymbol{r}_e[t]/k$ yields a Fourier transform analogous to Eq. (7).

For an investigation of the detectability of the quantum radiation in Eq. (7), the two-photon amplitude $|\mathfrak{A}_{\boldsymbol{k},\lambda,\boldsymbol{k}',\lambda'}|$

must be compared with the amplitude for the competing classical process, which is given by $|\alpha_{\boldsymbol{k},\lambda}\alpha_{\boldsymbol{k}',\lambda'}|$ to lowest order in $\alpha_{\boldsymbol{k},\lambda}$. In view of the smallness of the coupling g and assuming comparable results of the Fourier integrals (no resonances, etc.), there are basically two possibilities for achieving $|\mathfrak{A}_{\boldsymbol{k},\lambda,\boldsymbol{k}',\lambda'}| > |\alpha_{\boldsymbol{k},\lambda}\alpha_{\boldsymbol{k}',\lambda'}|$: small velocities $\dot{\boldsymbol{r}}_e^2 \ll 1$ or small angles ϑ between $\boldsymbol{k}$ and $\dot{\boldsymbol{r}}_e$ (blind spot). The first alternative is probably impractical since the total effect becomes too small, but the latter option can be realized for a unidirectional acceleration leading to a well-defined blind spot (cf. Fig. 1).

Ultrarelativistic regime.—In order to be as close as possible to the scenario of the Unruh effect (eternally uniform acceleration), let us consider an external electric field which is approximately constant over a sufficiently long period of time. Of course, such a pulse $E(t)$ necessarily accelerates the electrons to ultrarelativistic velocities $\gamma \gg 1$. Since both quantum and classical radiation will be boosted forward in this case (cf. Fig. 1), we shall focus on a small forward cone $\vartheta \ll 1/\gamma$; see also Eq. (12) below. In this limit, Eq. (7) simplifies to (for $\lambda = \lambda'$)

$$\mathfrak{A}_{\boldsymbol{k},\lambda,\boldsymbol{k}',\lambda'} \approx \int d\tau \frac{g\gamma(\tau)}{iV\sqrt{kk'}} e^{i(k+k')\tau}, \tag{9}$$

with a time-dependent Lorentz factor $\gamma(\tau)$ whose evolution is determined by $d\gamma/dt \approx qE(t)/m$ as well as $dt \approx 2\gamma^2 d\tau$. Since $\gamma(\tau)$ jumps from its initial value of order one to its maximum value $\gamma_{\max} \gg 1$ on an effective time scale $\Delta\tau$ much shorter than the pulse length Δt, the above integral behaves like the Fourier transform of a Heaviside step function for small k

$$\mathfrak{A}_{\boldsymbol{k},\lambda,\boldsymbol{k}',\lambda'} = \mathcal{O}\left(\frac{g\gamma_{\max}}{V\sqrt{kk'}(k+k')}\right). \tag{10}$$

The behavior at large k is determined by the structure of the integrand in (9) near the maximum Lorentz factor $\gamma_{\max}$,

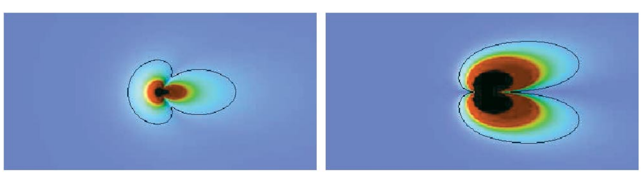

FIG. 1 (color). Two-photon amplitudes for quantum (Unruh, left image) and classical (Larmor, right image) radiation generated by an electron which is (after being initially at rest) accelerated by a Gaussian electric field pulse with a width of 0.3 as to a moderately relativistic velocity $\gamma_{\max} \approx 2$ and moves to the right. The amplitudes are shown for $\boldsymbol{k} = \boldsymbol{k}'$ and $\lambda = \lambda'$ and plotted as a function of $\boldsymbol{k}$. The points in the middle of the images correspond to $\boldsymbol{k} = 0$ and the maximum k values at the left and right boundaries to 30 keV. The color coding (same in both images) is chosen such that red indicates large amplitude and blue vanishing amplitude. The black areas are the excised singularities at $\boldsymbol{k} = 0$ and the black lines are isolines. One can clearly see that quantum radiation (left image) dominates inside a small forward and backward cone.

where $d\gamma/d\tau$ drops from $2\gamma^2 qE/m$ to zero on an effective time scale of $\Delta\tau = \mathcal{O}(\Delta t/\gamma_{\max}^2)$. As one would expect from the simplified picture of a Lorentz boosted Unruh temperature, the photons with the highest energies are typically created in the final stage of the acceleration phase. In order to resolve deviations (smoothening) from the step-function behavior, the wave number must exceed a certain cutoff $k_{\text{cut}} = \mathcal{O}(\gamma_{\max}^2/\Delta t)$ determined by the effective time scale $\Delta\tau = \mathcal{O}(\Delta t/\gamma_{\max}^2)$. Since $\gamma_{\max}$ is roughly given by the pulse length Δt times the acceleration qE/m, the cutoff wave number can alternatively be written as $k_{\text{cut}} = \mathcal{O}(\gamma_{\max} qE/m)$, which corresponds to the Unruh temperature (1) boosted by $\gamma_{\max}$. Below this cutoff, the amplitude behaves as (10), and, for k values larger than k_{cut}, the decline is faster than polynomial [assuming a smooth C^∞ pulse $E(t)$]. This behavior has also been confirmed by numerical integrations of Eq. (7). Hence, the typical photon energies are determined by that cutoff $k_{\text{cut}} = \mathcal{O}(\gamma_{\max} qE/m)$ and might even exceed the electron's rest mass for large Lorentz factors.

An analogous estimate for the Larmor radiation yields

$$\alpha_{\boldsymbol{k},\lambda} \approx \int d\tau \frac{2q\vartheta\gamma^2}{\sqrt{2Vk}} e^{ik\tau} = \mathcal{O}\left(\frac{q\vartheta\gamma_{\max}^2}{\sqrt{Vk^3}}\right), \tag{11}$$

with basically the same wave number cutoff. Of course, quantum radiation dominates for sufficiently small ϑ. For the cutoff wave number k_{cut}, the angular size of the small forward cone of "quantum domination" scales as

$$\vartheta_{\max} = \mathcal{O}\left(\sqrt{\frac{qE}{m^2}\frac{1}{\gamma_{\max}}}\right) \tag{12}$$

and is determined by $\gamma_{\max}$ and ratio of the electric field E over the Schwinger [8] limit $E_S = m^2/q$. The probability of (two-photon) emission in this cone is given by

$$\mathfrak{P}(\vartheta_{\max}) = \sum_{\boldsymbol{k},\lambda,\boldsymbol{k}',\lambda'}^{\vartheta<\vartheta_{\max}} |\mathfrak{A}^2_{\boldsymbol{k},\lambda,\boldsymbol{k}',\lambda'}| = \mathcal{O}\left(q^4 \frac{E^4}{E_S^4}\right). \tag{13}$$

Interestingly, for a given electric field strength E, this probability does not depend significantly on the pulse length since $\gamma_{\max}$ cancels—but the energy $k_{\max}$ and the angular distribution $\vartheta_{\max}$ of the emitted photons does.

Extensions.—Since the observation of the photon pairs requires electromagnetic fields which are not too far below the Schwinger limit, one should also consider the impact of these ultrahigh fields on the QED vacuum, which then acts as a medium and displays effects such as birefringence. To this end, we consider the first nonlinear corrections from the Euler-Heisenberg Lagrangian [11]

$$\mathcal{L} = \frac{1}{2}(\boldsymbol{E}^2 - \boldsymbol{B}^2) + \frac{(\boldsymbol{E}^2 - \boldsymbol{B}^2)^2 + 7(\boldsymbol{E}\cdot\boldsymbol{B})^2}{90\pi E_S^2/\alpha_{\text{QED}}}. \tag{14}$$

If we neglect the external magnetic field and linearize around an approximately homogeneous external electric field $\boldsymbol{E}_0$, we get $\boldsymbol{D} = \underline{\boldsymbol{\varepsilon}} \cdot \boldsymbol{E}$, with the permittivity tensor $\underline{\boldsymbol{\varepsilon}}$.

Hence, the transversality condition $\boldsymbol{k}\cdot\underline{\boldsymbol{\varepsilon}}\cdot\boldsymbol{e}_{\boldsymbol{k},\lambda}=0$ deviates from $\boldsymbol{k}\cdot\boldsymbol{e}_{\boldsymbol{k},\lambda}=0$, and, thus, the Larmor radiation does not necessarily vanish in the forward direction anymore. However, for a trajectory along a field line $\dot{\boldsymbol{r}}_e \parallel \boldsymbol{E}_0$, i.e., along an eigenvector of $\underline{\boldsymbol{\varepsilon}}$, the Larmor radiation still has the blind spot $\boldsymbol{k}\perp\boldsymbol{e}_{\boldsymbol{k},\lambda}$ in the forward direction $\boldsymbol{k}\parallel\boldsymbol{E}_0$.

With an external magnetic field $\boldsymbol{B}_0$, on the other hand, additional terms appear. Assuming $\boldsymbol{k}\parallel\boldsymbol{E}_0\perp\boldsymbol{B}_0$, one of the (linear) polarizations λ has a blind spot in the forward direction $\boldsymbol{k}\cdot\boldsymbol{e}_{\boldsymbol{k},\lambda}=0$ but the other one λ' has not $[\boldsymbol{k}\cdot\boldsymbol{e}_{\boldsymbol{k},\lambda'}=4\alpha_{\rm QED}kE_0B_0/(45\pi E_S^2)]$. Although the numerical prefactor is rather small (2×10^{-4}), this effect should be taken into account when searching for quantum radiation [12]. On the other hand, it might also provide an opportunity for testing the birefringent nature of the QED vacuum in the presence of ultrahigh external fields (see also [13]). Similarly, one obtains small corrections to the two-photon correlations.

Summary.—We have studied the conversion of (virtual) quantum vacuum fluctuations into (real) particles by noninertial (planar) Thomson scattering for the example of strongly accelerated electrons. This quantum radiation can be discriminated from classical (Larmor) radiation via the different angular (blind spot cf. Fig. 1) and spectral distributions and the distinct photon statistics: In the quantum case, the photons are always emitted in pairs (squeezed state) with perfectly correlated polarizations—whereas the classical (Larmor) radiation corresponds to a coherent state (independent photons with Poissonian statistics).

The probability of emitting two photons with wave numbers $k\ll k_{\rm cut}$ scales as $1/k^4$ for quantum (Unruh) radiation and as $1/k^6$ for Larmor. Note that the spectrum is not Planckian, in general. This is caused by the nontrivial translation from the accelerated frame to the inertial frame with a time-dependent Lorentz boost factor $\gamma(t)$ whose rate of change is of the same order as the frequency corresponding to the Unruh temperature (i.e., nonadiabatic). In addition, a state consisting of pairs of correlated photons can never be exactly thermal.

Apart from the derivation of the correct two-photon spectral and angular distributions and the perfect correlation of polarization, the approach presented in this Letter (properly accounting for the two-photon nature of quantum radiation) presents another advantage in comparison with previous works [7]: Because the probability for a single Larmor photon in the cone (12) scales as $q^2E^2/E_S^2\gg\mathfrak{P}(\vartheta_{\rm max})$, coincidence measurements are probably crucial for discriminating classical from quantum radiation (coherent versus squeezed state).

According to Eq. (13), the signatures of the Unruh effect might be detectable in future facilities (cf. [14–16]) generating $\mathcal{O}(1\ {\rm as})$ electric field pulses not too far below [9] the Schwinger limit $E_S\approx1.3\times10^{18}$ V/m which accelerate the electrons (e.g., created via the Schwinger effect) to ultrarelativistic velocities. The emitted photons could be measured via their Compton scattering in large-volume Si(Li) orthogonal-strip detectors (placed in the blind spot), allowing one to determine their wave number $\boldsymbol{k}$ and polarization λ (to deduce the photon entanglement). Conversely, one might use the spectral and angular distribution of the radiation to infer the characteristic parameters of the pulse such as E, Δt, and $\gamma_{\rm max}$, etc.

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