

Shallow-angle inverse Compton scattering: Experimental demonstration and future applications

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Inverse Compton scattering (ICS) at shallow angles can be used to tune the spectrum and increase the brightness of the emitted radiation. Here, we investigate the geometric scaling laws governing this interaction and experimentally validate them by demonstrating shallow-angle ICS of a 780 nm laser crossing a 4.7 MeV electron beam at 5.8°. Angular and spectral measurements of the emitted radiation show excellent agreement with analytical and numerical predictions. To further optimize this interaction, we consider the use of velocity-matched pulse-front tilt to decouple the total yield from the interaction bandwidth. The resulting framework is used to evaluate the performance scaling of three distinct driving accelerator architectures: a radio-frequency linac, a laser wakefield accelerator, and a storage ring synchrotron. These results establish a scalable, practical pathway for deploying next-generation, high-brightness compact light sources across diverse accelerator platforms.

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I. INTRODUCTION

X-ray sources permit the interrogation of a wide range of materials, including integrated circuits, cells, and electrocatalysts [1–3]. Unfortunately, high-brightness sources in the x-ray spectral range are limited to large scale facilities such as synchrotrons and free electron lasers [4–6]. Further routes toward compact and powerful x-ray sources are therefore a subject of active research. A promising approach is inverse Compton scattering (ICS) radiation from relativistic electrons colliding with a laser pulse [7–12]. The term inverse here indicates that the scattered laser photons gain energy from the electrons. In the standard head-on geometry, where the laser pulse counterpropagates with respect to the electron beam, the Doppler upshift of the scattered photons is maximized, allowing x-ray photon energies to be produced with comparatively modest electron beam energies. This feature is a major attraction of ICS-based x-ray generation schemes, because it enables compact, small-scale accelerator implementations [13–16]. The drawback is that the low electron energy also corresponds to a relatively small Lorentz factor, producing a large characteristic emission angle for the radiation that scales approximately

as $(1/\gamma)$ [17], which together with the small cross section [18] significantly restricts the source brightness compared to large scale facilities.

The crossing angle θ_0 of the laser pulse with the e -beam allows for additional tuning of the radiation wavelength, which is given by [19]

$$\lambda_{\text{rad}} = \frac{\lambda_0}{2\gamma^2(1 - \beta \cos \theta_0)} \left(1 + \gamma^2 \theta^2 + \frac{a_0^2}{2} \right), \quad (1)$$

where λ_0 is the wavelength of the laser beam, β is the velocity of the electrons normalized to the speed of light c , θ is the scattering angle from the e -beam axis, and a_0 is the normalized vector potential of the laser beam. Here and in the following, we assume that the electron beam is relativistic ($\gamma \gg 1$) and redshifting from ponderomotive forces is negligible ($a_0 \ll 1$). Equation (1) indicates that the radiation wavelength is tunable via the crossing angle from the visible to the hard x-ray spectral range for MeV-scale beams (see Fig. 1). Equivalently, the same x-ray wavelength can be produced with a higher electron beam energy by moving away from the head-on geometry and reducing the laser-electron crossing angle. For example, with a 1 μm laser, radiation at $\lambda_{\text{rad}} = 2.2$ nm, corresponding to a photon energy of 0.6 keV, can be produced either with a 5 MeV electron beam in the conventional head-on geometry, ($\theta_0 = 180^\circ$), or in a shallow near co-propagating geometry, ($\theta_0 \simeq 6.3^\circ$) using a 100 MeV electron beam, still within the range of compact accelerator setups [20,21].

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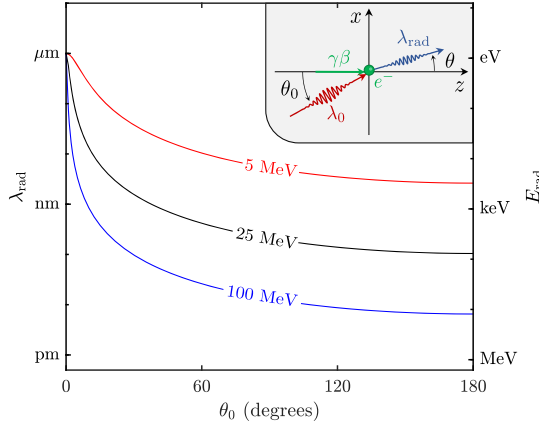


FIG. 1. On-axis radiation wavelength and photon energy tuning by crossing angle for three electron beam energies (indicated by color) assuming a laser wavelength of 1 μm . Inset: schematic of Compton scattering with driving laser in red, electron in green, and scattered radiation in blue.

The use of the higher electron beam energy immediately improves brightness through relativistic beaming as the characteristic emission cone narrows approximately as $(1/\gamma)$, concentrating the photons into a smaller solid angle. In addition, there are several other advantages of employing a shallow-angle scattering geometry. At the higher e -beam energy, space charge forces and geometric emittance are reduced significantly allowing for tighter transverse focusing and improved compression leading to even brighter radiation pulses [22]. Second, the near co-propagating laser can act as an ultrafast temporal gate on the electron bunch, enabling ultrashort x-ray pulse generation without requiring the full electron bunch to be compressed to the same duration; in this regime, the radiation pulse length can be limited primarily by the laser pulse duration [23,24]. Third, in combination with a laser pulse-front tilt, a shallow crossing angle can avoid restrictions on the photon yield imposed by laser diffraction and spectral broadening due to high laser intensities [25–28]. Finally, the higher e -beam energy permits the use of inverse free electron laser bunching [29–31], which can be used to impose density modulations down to the nanometer-scale [32], unfolding the potential of super-radiant Compton scattering [33,34].

Based on this, a multitude of ICS sources [19,34,35], including laser driven free electron lasers [26,27], have been proposed in recent years involving such an overtaking scattering geometry. Despite the wide interest, the properties of shallow-angle ICS have remained largely unexplored: the smallest crossing angle studied experimentally is 90° [9,36,37] and recently 30° [38]. Here, we provide a comprehensive study of the shallow-angle ICS regime, combining theoretical modeling, experimental demonstration, and machine design applications. First, we establish a simple but robust analytical framework quantifying the

geometric dependencies of the photon yield, interaction length, and spectral brightness in an overtaking geometry. We then validate these scaling laws experimentally by reporting the observation of inverse Compton scattering from a 4.7 MeV electron beam crossing a 780 nm laser at a shallow angle of 5.8° . The visible spectral range of the emitted radiation allows us to use high-sensitivity detectors to characterize the interaction in detail. The measured angular, spectral, and yield scalings show excellent agreement with both the analytical model and start-to-end Liénard-Wiechert simulations, thereby benchmarking the framework in a regime where the radiation can be extensively diagnosed. Having validated the model experimentally, we then extend it to optimized interaction geometries, including velocity-matched pulse-front tilt, and apply it to evaluate distinct electron-driver architectures for future high-brightness shallow-angle ICS light sources.

II. PROPERTIES OF SHALLOW-ANGLE INVERSE COMPTON SCATTERING

The spatiotemporal overlap in an overtaking geometry differs strongly from conventional head-on scattering, modifying radiation properties including the photon yield, linewidth, divergence, and, consequently, the source brightness. In this section, we establish the analytical framework governing these properties, concentrating the discussion on shallow interaction angles. For a more general discussion on arbitrary crossing angles and brightness contributions, we refer the reader to [8,39], respectively.

For the following calculations, we define the coordinate system such that the electron beam propagates along the z axis with velocity βc and the laser pulse crosses the electron beam in the x - z plane at an incidence angle θ_0 , where $\theta_0 = 0$ corresponds to co-propagating beams. The electron beam is modeled as a 3D Gaussian density distribution with N_e electrons and rms spatial sizes $\sigma_{x,y,z}$ given by $\rho_e = N_e / [(2\pi)^{3/2} \sigma_x \sigma_y \sigma_z] \exp[-x^2/(2\sigma_x^2) - y^2/(2\sigma_y^2) - (z - \beta ct)^2/(2\sigma_z^2)]$. Similarly, the photon density for N_l laser photons is described by a 3D Gaussian distribution with out-of-plane (s) and in-plane (p) transverse rms waists $\sigma_{s,p}$ and rms laser pulse length σ_t given by $\rho_l = N_l / [(2\pi)^{3/2} \sigma_p \sigma_s c \sigma_t] \exp[-s^2/(2\sigma_s^2) - p^2/(2\sigma_p^2) - (l - ct)^2/(2c^2\sigma_t^2)]$. Here, $s = y$, $p = x \cos \theta_0 - z \sin \theta_0$, and $l = x \sin \theta_0 + z \cos \theta_0$ are the transverse and longitudinal coordinates of the laser pulse. The rms intensity waists are related to the $1/e^2$ field waists by $\sigma_{s,p} = w_{s,p}/2$.

A. Photon yield

The total scattered photon yield from a single interaction can be written as $N_{\text{ph}} = \sigma_T \mathcal{L}$, where σ_T is the Thomson cross section and $\mathcal{L}$ is the time-integrated luminosity. The luminosity is obtained by integrating the product of

the beam densities over all space and time, scaled by the Møller relative velocity factor $K = c(1 - \beta \cos \theta_0)$, yielding $\mathcal{L} = c(1 - \beta \cos \theta_0) \int_{-\infty}^{\infty} dt d^3r \rho_e(\mathbf{r}, t) \rho_l(\mathbf{r}, t)$. Evaluating this multivariate Gaussian overlap integral provides a general expression for the photon yield at an arbitrary crossing angle:

$$N_{\text{ph}} = \frac{\sigma_T N_e N_l}{2\pi \Sigma_s \Sigma_p}, \quad (2)$$

where Σ_s and Σ_p represent the effective out-of-plane and in-plane interaction widths, respectively. Note that this expression is valid only in the linear regime where $a_0 \ll 1$. The out-of-plane interaction width is simply the quadrature sum of the transverse sizes, $\Sigma_s = (\sigma_s^2 + \sigma_y^2)^{1/2}$, which is identical to the head-on case. The in-plane interaction width Σ_p couples the transverse and longitudinal spot sizes to the crossing angle:

$$\Sigma_p = \left[\sigma_p^2 + \sigma_x^2 \left(\frac{\beta - \cos \theta_0}{1 - \beta \cos \theta_0} \right)^2 + \frac{(\sigma_z^2 + \beta^2 c^2 \sigma_t^2) \sin^2 \theta_0}{(1 - \beta \cos \theta_0)^2} \right]^{1/2}. \quad (3)$$

This expression explicitly captures the geometric degradation of luminosity. In a perfectly head-on collision, the longitudinal terms vanish entirely, maximizing the yield via the purely transverse overlap. Conversely, at shallow angles, the longitudinal terms typically dominate the interaction volume, making longitudinal compression of the beams critical for maximizing the yield. This becomes apparent by expanding the angular projection term $\sin \theta_0 / (1 - \beta \cos \theta_0) \simeq 2\gamma^2 \theta_0 / (1 + \gamma^2 \theta_0^2)$ for small angles and highly relativistic electron beams. This geometric penalty maximizes to γ at the crossing angle $\gamma\theta_0 = 1$, and asymptotically approaches $2/\theta_0$ for macroscopic shallow angles ($\gamma\theta_0 \gg 1$). Notably, at $\gamma\theta_0 = 1$, where the projected axial velocity of the laser matches the electron beam velocity, the geometric penalty from the transverse electron bunch size σ_x vanishes entirely, although in this case the Compton edge is restricted to frequency doubling.

To demonstrate the sensitivity to bunch length, Fig. 2 plots the photon yield N_{ph} and the angular flux $\gamma^2 N_{\text{ph}}$ as a function of the electron beam energy. For this comparison, the ICS source is continuously tuned to emit 50 keV photons by actively adjusting the interaction angle θ_0 as the beam energy increases. Because higher energies require increasingly shallow crossing angles to maintain the 50 keV resonance, the photon yield degrades across all cases, suffering most severely for the longest longitudinal bunch distributions. However, the angular flux increases significantly with beam energy, particularly for the 10 fs case, explicitly demonstrating that longitudinal compression of the electron beam is imperative for maximizing brightness in shallow-angle Compton scattering sources.

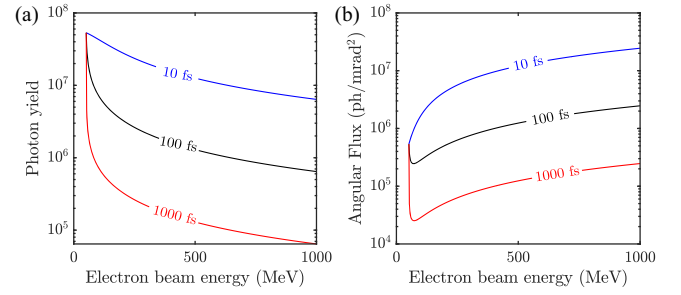


FIG. 2. (a) Photon yield and (b) angular flux plotted against the electron beam energy for an ICS source continuously tuned to a 50 keV scattered photon energy via the interaction angle. Parameters are $\lambda_0 = 1 \mu\text{m}$ with 100 mJ laser pulse energy and 100 pC bunch charge, $\sigma_{x,y} = \sigma_{s,p} = 10 \mu\text{m}$, for three distinct temporal durations σ_t distinguished by color.

Beyond these geometric penalties, the crossing angle also influences the microscopic emission dynamics. This effect is most pronounced at $\gamma\theta_0 \simeq 1$, corresponding to the relativistic Brewster angle, where the laser intersects the electrons at exactly 90° in their instantaneous rest frame. This kinematic condition drives a strong polarization-dependent response; for p -polarized laser pulses, the induced dipole aligns directly with the electron beam axis, fully suppressing on-axis emission in close analogy with Brewster reflection at dielectric interfaces [40]. Even as the interaction is scaled to larger shallow angles ($\gamma\theta_0 \gg 1$) to access shorter wavelengths, polarization continues to influence the radiation profile by introducing a transverse asymmetry that shifts the maximum intensity slightly off axis. Crucially, while the laser polarization state dictates the exact angular emission pattern of each electron, it remains completely decoupled from the macroscopic spatial overlap of the beams. Therefore, the geometric yield scaling derived here serves as the universal baseline for predicting the total photon yield across all crossing angles.

B. Interaction length

The spectral and angular characteristics of the radiation are strongly dictated by the interaction length, i.e., the number of optical periods a single electron traverses during the collision. By substituting the nominal single-electron trajectory $z = c\beta t$ and $x = y = 0$ into the laser photon distribution ρ_l , one can readily show that the interaction length N_φ (expressed in radians) is given by

$$N_\varphi = \omega_0 \sigma_t \left[1 + \frac{c^2 \sigma_t^2 \beta^2 \sin^2 \theta_0}{\sigma_p^2 (1 - \beta \cos \theta_0)^2} \right]^{-1/2}. \quad (4)$$

The bracketed term represents a direct competition between the longitudinal transit time $t_l = \sigma_t / (1 - \beta \cos \theta_0)$ and the transverse transit time across the in-plane waist $t_p = \sigma_p / (c\beta \sin \theta_0)$. Consequently, the interaction length

can be actively tuned via the laser waist σ_p , but ultimately remains bounded by the physical laser pulse length. Thus, the interaction length is typically shorter than the head-on equivalent unless ultrashort pulses are used such that $t_l \ll t_p$. The limitation imposed by the laser pulse length can be systematically circumvented by applying an appropriate spatiotemporal pulse-front tilt to the laser as discussed later.

Crucially, N_ϕ dictates the emission duration for a single electron, thereby establishing the intrinsic spectral bandwidth as $\Delta\omega/\omega = 1/N_\phi$. This represents the theoretical minimum spectral bandwidth, before accounting for inhomogeneous broadening effects driven by the macroscopic electron beam's finite emittance and energy spread. Due to the coupling of frequency and emission angle via the Doppler shift [Eq. (1)], the opening angle of the central on-axis radiation is also determined by the intrinsic bandwidth, scaling as $\sigma_{\theta,\text{rad}} = (\Delta\omega/\omega)^{1/2}/\gamma = (N_\phi)^{-1/2}/\gamma$. A longer interaction length thus produces a narrower bandwidth and increased collimation of the central wavelength, and therefore leads to a brighter radiation source.

C. Sensitivity to spatiotemporal alignment

One advantage of shallow-angle Compton scattering is that the elongated overlap geometry relaxes transverse alignment tolerances in the scattering plane. A spatial offset in this direction primarily shifts the effective collision point along the beamline rather than preventing the laser and electron beams from intersecting. Provided this shift remains within the laser depth of focus, the luminosity is only weakly affected. The tradeoff is a much tighter temporal tolerance. Because the laser pulse and electron bunch have a reduced relative longitudinal velocity in an overtaking geometry, a time-of-arrival error directly reduces their envelope overlap and suppresses the photon yield [23]. This contrasts with the head-on configuration, where the beams pass through one another rapidly and small timing errors produce a weaker reduction in overlap. This geometric penalty is particularly severe for the ultrashort bunches required to maximize the photon yield. We can analytically quantify this degradation by introducing a time-of-arrival error Δt , which results in a Gaussian suppression of the single-shot yield:

$$N_{\text{ph}}(\Delta t) = N_0 \exp \left[-\frac{\beta^2 c^2 \Delta t^2 \sin^2 \theta_0}{2 \Sigma_p^2 (1 - \beta \cos \theta_0)^2} \right], \quad (5)$$

where $N_0 = N_{\text{ph}}(0)$ is the peak yield for a perfectly synchronized collision. In an experimental setup, the timing error fluctuates from shot to shot, which can be accounted for by convolving this expression with the time-of-arrival jitter distribution. Because the in-plane interaction width Σ_p is typically heavily dominated

by the longitudinal components at shallow angles, the trigonometric dependencies almost entirely cancel out. The argument of the exponent roughly simplifies to $-\Delta t^2/[2((\sigma_z/\beta c)^2 + \sigma_t^2)]$. Consequently, any time-of-arrival drift directly suppresses the photon yield and brightness, necessitating relative synchronization control strictly bounded by the combined scale of the laser and electron bunch lengths.

III. EXPERIMENTAL SETUP AND CALIBRATION

To experimentally validate these theoretical scaling laws and demonstrate the physical feasibility of the shallow-angle regime, a dedicated ICS experiment was designed and executed. The experiment was carried out at the UCLA Pegasus photoinjector [41]. A schematic of the experimental geometry showing the shallow angle interaction and the diagnostic used to spectrally and angularly resolve the visible light radiation is shown in Fig. 3. A 200 pC electron bunch is emitted by illuminating a high quantum efficiency NaSkB photocathode [42] in a S-band 1.6 cell rf gun with a frequency tripled Ti:Sa laser system providing 260 nm 100 fs FWHM UV pulses. The gun is operated at 20° phase and acceleration gradient of 65 MV/m and is followed by a buncher linac tuned at -80° from crest with a gradient of 17 MV/m to compress the 4.7 MeV beam to 300 fs rms at the interaction point (IP) with the laser. A radially magnetized permanent magnet (PM) solenoid [43] located 8 cm upstream of the IP strongly focuses the electrons to a waist of about 150 μm at the interaction point. A PM dipole is used to separate the e -beam from the Compton radiation and minimize the radiation noise on the camera.

The infrared (IR) line that delivers the 780 nm laser pulse with 6 mJ pulse energy and 100 fs FWHM pulse length at the interaction point is first split by a polarizing beam splitter (PBS) from the same laser line that drives the photoinjector. Downstream of the PBS, the polarization angle ϕ_0 of the linear polarized pulse is controlled by a half wave plate, which is kept normal to the interaction plane $\phi_0 = 0$ (s-polarized). The laser pulse is then focused by a lens with a focal length of $f = 250$ mm to a $1/e^2$ waist of 100 μm in the IP corresponding to a peak intensity of 3.5×10^{14} W/cm². Upstream the lens we inserted a

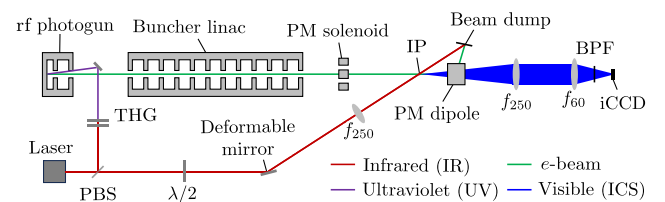


FIG. 3. Schematic of the beamlines used in the experiment. See text for details.

deformable mirror to compensate for astigmatism induced in the transport of the laser beam and that allowed live-adjustment to the laser spatial profile at the IP.

The Compton diagnostic line is a two lens point source microscope with a demagnification of 24% to increase sensitivity to low photon counts collected on the pixels of 12.8 μm of the intensified (i)CCD detector (Princeton Instruments PI-MAX 4). The first lens after the interaction is a 2 in. diameter 30 cm focal length lens which defines the maximum collection angle to 0.09 rads. The use of shorter focal length lens to maximize the collection angle was not possible due to the space to accommodate a large gap permanent magnet dipole and eliminate the radiation noise from the electron beam hitting the optics. To filter out any background laser light, we use a strong (OD > 15) band pass filter (BPF) allowing only radiation to pass in the visible (400–720 nm).

Absolute photon yield calibration of the diagnostic line was performed using forward optical transition radiation (TR) from a metal foil at the IP. TR provides an ideal calibration source due to its well-defined analytical properties and similarity to the Compton radiation under study. The spectral angular photon distribution of TR from an electron crossing the metal-vacuum interface can be written as [44]: $d^2N_{\text{ph}}^{\text{TR}}/(d\omega d\Omega) = \mu_0 c e^2 / (4\pi^3 \hbar \omega) \beta^2 \sin^2 \theta / (1 - \beta^2 \cos^2 \theta)^2$. The collected TR yield per electron is given by $N_{\text{ph}}^{\text{TR}} = 2\pi \int_0^{\theta_c} \sin \theta d\theta \times \int_0^\infty SS(\omega) d^2N_{\text{ph}}^{\text{TR}}/(d\omega d\Omega) d\omega$, where SS is the spectral detection sensitivity of the Compton diagnostic line given by the product of the quantum efficiency (QE) of the camera and the transmission efficiency (TE) of the optics as shown in the bottom inset of Fig. 4(a). For the maximum collection angle of $\theta_c = 0.09$ rad, and a 6.4 MeV electron beam used for the calibration, the detected number of photons per electron is $N_{\text{ph}}^{\text{TR}} = 1.5 \times 10^{-4}$. The absolute system response was determined to be ~ 1 photon/count at

a linear detector gain of 1000, establishing a highly sensitive baseline for the ICS photon counting.

As the spatiotemporal window for Compton scattering is very small, we use intermediate diagnostics to assist in alignment. For the spatial overlap, a removable lanex screen was used, allowing for transverse profiling of both the laser and e -beam at the IP. Temporal overlap was then established via electro-optic sampling (EOS) of the electron bunch's Coulomb field [45,46]. A birefringent ZnTe crystal, placed at the IP between crossed polarizers, measured the field-induced birefringence from the traversing electron beam as probed by the strongly attenuated IR laser pulse. The resulting polarization retardation angle is given by $\Gamma = 2\pi L n_0^3 r_{41} T_{\text{ZnTe}} E_C / \lambda_0$ [47], where $L = 1$ mm is the crystal length, $n_0 = 2.85$ is the refractive index of ZnTe at $\lambda_0 = 780$ nm, $r_{41} = 4 \times 10^{-12}$ m/V is the EO coefficient of ZnTe, and $T_{\text{ZnTe}} \approx 0.5$ is the correction coefficient accounting for Fresnel losses, absorption, and dephasing effects [48]. This retardation translates into a measurable intensity modulation governed by Malus' law. By scanning the relative arrival time using a linear delay stage, we obtain a time trace of the EOS signal [Fig. 4(b)], allowing the temporal overlap to be found with subpicosecond accuracy.

The EOS trace is also utilized to characterize the time-of-arrival (TOA) jitter between the electron bunch and the laser pulse. Specifically, the TOA jitter is assumed to be Gaussian distributed with standard deviation that is determined through an iterative fitting procedure: We apply a Wiener deconvolution (see Appendix B for details) to the raw EOS trace for a range of rms jitter values and select the value of σ that minimizes the spread of the signal counts relative to the average at each time step. The extracted TOA rms jitter of 400 fs is subsequently used to deconvolve the measured Compton scattering traces. In addition, the resulting deconvolved profile, shown in Fig. 4(b), is used to determine the electron bunch length and peak current.

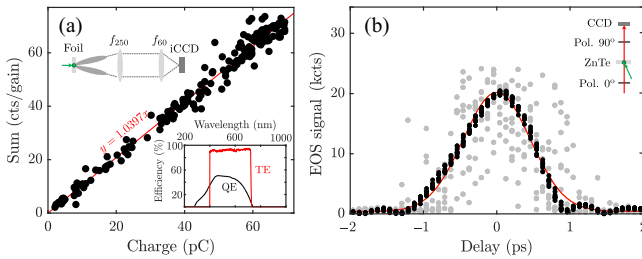


FIG. 4. (a) Transition radiation counts (normalized to detector gain) against bunch charge used to calibrate iCCD. The red line is a linear fit through the data. Top inset: schematic of the experiment. Bottom inset: quantum efficiency and transmission efficiency of the diagnostic line. (b) EOS signal (gray dots) and Wiener deconvolved EOS signal (black dots) as a function of IR delay from a 200 pC e -bunch. The red curve is a Gaussian fit through the deconvolved data. Inset: schematic of the EOS line.

IV. START-TO-END SIMULATIONS

The measured signals are compared to simulations starting from the photoemission of the electrons at the cathode to the shallow-angle Compton radiation by collision with the laser pulse. The numerical calculations consist of two parts: First, particle tracking simulations are performed with the general particle tracer (GPT) code [49]. GPT integrates the equations of motion of the electrons in electromagnetic fields without approximations, including self-forces such as space charge. Field maps of the accelerator structures and permanent magnet solenoid are used to simulate the transport of the electron beam. The laser pulse is modeled by a linearly polarized Gaussian pulse given by $\mathbf{E}_{\text{laser}} = E_0 u_s u_p \exp[-(l - ct)^2 / (4\sigma_l)] \mathbf{e}$, where $u_\xi = \text{Re}\{q_\xi^{-1/2} \exp[-ik_0 \xi^2 / (2q_\xi)]\}$ the spatial

dependence in the direction with $q_\xi = l + iz_{R,\xi}$ and $z_{R,\xi} = \pi w_\xi^2/\lambda_0$ the Rayleigh ranges for $\xi \in \{s, p\}$ and $\epsilon = \cos \theta_0 \sin \phi_0 \mathbf{e}_x + \cos \phi_0 \mathbf{e}_y - \sin \theta_0 \sin \phi_0 \mathbf{e}_z$ the polarization vector.

Second, the trajectories of the particle tracking simulations, in particular from the interaction with the laser pulse, are used to calculate the radiation field in post-processing using the Liénard-Wiechert fields [50]: $\mathbf{E}_{\text{rad}}(\mathbf{r}, t) = q/(4\pi\epsilon_0 c)[\mathbf{n} \times \{(\mathbf{n} - \boldsymbol{\beta}) \times \dot{\boldsymbol{\beta}}\}/(1 - \mathbf{n} \cdot \boldsymbol{\beta})^3/R]_{\text{ret}}$ with $\mathbf{n}$ is the unit vector connecting the electron to the observer positioned at $\mathbf{r}$, $\boldsymbol{\beta}$ is the normalized velocity vector, $\dot{\boldsymbol{\beta}}$ is the normalized acceleration vector and $R = |\mathbf{r} - \mathbf{r}_e(t)|$ with $\mathbf{r}_e(t)$ is the trajectory of an electron and $\mathbf{r}$ the observer position. The term between square brackets is evaluated at the retarded time $t_r = t - R/c$, with t the observer time. Since the simulation (retarded) time is discrete we interpolate to calculate the field at specified observer times, which allows for summing of the fields from each simulated particle such that $\mathbf{E}_{\text{rad}}(\mathbf{r}, t) = \sum_j \mathbf{E}_{\text{rad},j}(\mathbf{r}, t)$ with $\mathbf{E}_{\text{rad},j}(\mathbf{r}, t)$ the field of the j th particle. In order to calculate the yield the fields are first Fourier transformed $\mathbf{E}_{\text{rad}}(\mathbf{r}, \omega) = \mathcal{F}(\mathbf{E}_{\text{rad}}(\mathbf{r}, t))$ and filtered with the spectral sensitivity of the collection system. Then, the resulting spectral angular photon distribution $d^2 N_{\text{ph}}/(d\Omega d\omega) = c\epsilon_0 |\mathbf{r} \mathbf{E}_{\text{rad}}(\mathbf{r}, \omega)/(\hbar\omega)|^2$ is integrated over the collected solid angle and angular frequency, taking into account the spectral sensitivity of the detection line.

A. Validation with beam measurements

The particle tracking simulations are initialized using the machine parameters recorded during the experiment and validated by matching the output to a set of separate e -beam diagnostic measurements. These are summarized alongside the schematic of the beamline in Fig. 5. Starting at the cathode, we launch a nearly uniform transverse beam distribution with a radius of about 0.5 mm, according to the laser spot size as measured on a virtual cathode, as shown in Fig. 5(a), with the known 100 fs FWHM pulse length of the UV laser pulse. The longitudinal phase space distribution out of the gun as reconstructed using longitudinal phase space tomography [51] matches well with the expectations from simulations as shown in Fig. 5(b). The buncher linac gradient is calibrated using the e -beam energy for different phases as measured by a downstream dipole spectrometer. This was done for a lower charge bunch (100 pC) close to crest so that the full energy spectrum fits in the field of view of the spectrometer allowing for accurate energy measurements. Figure 5(c) shows the good agreement between the measurements and calibrated simulations. At the interaction point, several diagnostics were present: the spot size of the electron beam (and laser beam) was measured on the lanex screen as a function of charge, and coincides well with the simulations

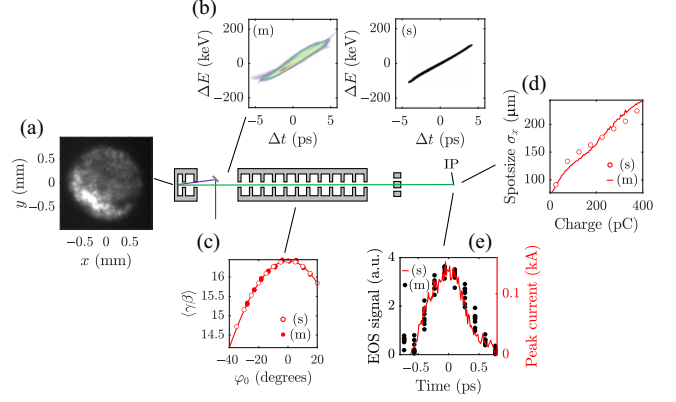


FIG. 5. Comparison of e -beam (m) measured observables with (s) particle tracking simulations: (a) UV spot size on cathode. (b) Longitudinal phase space distribution. (c) Average momentum as a function of input linac phase. (d) Spot size in interaction point. (e) EOS signal and simulated peak current.

as seen in Fig 5(d). Finally, the current distribution from the simulations agrees with the measured (and Wiener deconvolved) EOS time trace for the 100 pC case as shown in Fig. 5(e).

V. EXPERIMENTAL RESULTS

After overlapping in space and time the electron and laser beams we recorded on the iCCD the bright spot shown in Fig. 6(a) recording the spatial intensity distribution of the resulting Compton scattering. The rms source size is $88 \times 108 \mu\text{m}^2$, smaller than the measured e -beam size on the lanex and OTR screens, consistent with expectation from the convolution of the laser and electron distributions at the IP. The integrated yield captured within the maximal collection angle of $\gamma\theta_c = 0.8$ was typically 300 photons, which, taking into account the losses of the collection system, corresponds to about 3×10^3 emitted photons. This is in good agreement with the start-to-end simulations discussed in the previous section. Using

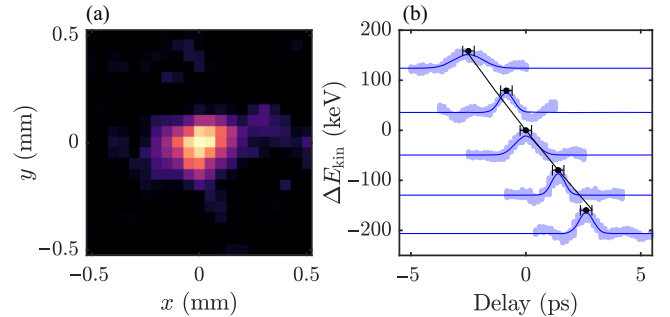


FIG. 6. (a) Intensity measured on the iCCD averaged over 85 shots. (b) Energy of the beam relative to the nominal 4.7 MeV as a function of arrival time at the IP represented by the black data points. The blue scattered data are the photon yield for each time delay.

additional low-pass filters with cutoff edges $\lambda_{\text{rad}} < 450$ nm and $\lambda_{\text{rad}} < 550$ nm, respectively, we verified the emission spectrum to be consistent with the calculated on-axis radiation wavelength $\lambda_{\text{rad}} = 414$ nm.

Shallow-angle Compton scattering exhibits severe sensitivity to time-of-arrival changes of the e -beam as discussed in Sec. II. We experimentally characterized this by systematically varying the electron bunch arrival time via its energy using the buncher linac and scanning the laser delay, as shown in Fig. 6(b). The Compton time traces, indicated by the blue scattered data, were deconvolved using a Wiener filter to remove the 400 fs TOA jitter of the rf compressed beam (see Appendix B). The centroids of the deconvolved traces map the energy-dependent arrival time, which agrees perfectly with the calculated drift $\Delta t = \Delta \gamma L / (c\beta^2 \gamma^3)$, with $L = 197$ cm the distance between the exit of the linac and the IP, as indicated by the solid black curve in the figure. Additionally, the rms duration of the Compton signal, found by fitting a Gaussian distribution to the time traces (solid blue curves), is consistent with the bunch length obtained from independent EOS measurements, further validating our analysis and confirming the prompt nature of the interaction.

The photon yield scaling of shallow-angle Compton scattering was then investigated. In general, the ICS photon yield by an electron bunch given by Eq. (2) scales linearly with the number of electrons N_e . However, the collected yield as a function of the charge, measured nondestructively with an integrating current transformer, is only linear up until about 150 pC as indicated by the black dots on Fig. 7(a). At higher charges, the yield starts to saturate as space charge forces oppose the compression of the e -beam and therefore increase the in-plane interaction width with the co-propagating laser pulse, which approximately scales as $\Sigma_p \propto \sigma_z / (c\sigma_t)$ since in our case the electron bunch length σ_z is much longer than the laser pulse length $c\sigma_t$. Consistently, the photon yield is observed

to scale with the peak current as shown by the red dots in Fig. 7(a), which was inferred independently using EOS. Note that this behavior is a unique feature of shallow-angle Compton not present in head-on scattering where Σ_p is independent of the bunch length as long as diffraction effects are negligible. The measurements here are in good agreement with the start-to-end numerical calculations, given by the white dots showing only a minor deviation post saturation, potentially caused by clipping due to space charge dominated beam blowup at, e.g., the edges of the PM solenoid.

In another set of measurements, we changed the laser pulse shape in the IP by imparting aberrations to the laser pulse with the deformable mirror. In particular, a nonzero astigmatism was applied to produce a wider spot in (p) and out (s) of the plane of interaction, independently, while the pulse energy remained constant. The insets in Fig. 7(b) show the intensity distributions of the laser at the IP as measured on the lanex screen. The resulting collected yield for each laser spot, in very good agreement with numerical simulations displayed in Fig 7(b), reveals the typical (for head-on scattering) inverse scaling with the waist w_s in the s direction, but shows an almost constant yield for varying waist w_p in the p direction, despite equivalent reduction in laser intensity. The asymmetry arises in a side-scattering geometry since the interaction length is affected by the waist w_p [34,52]. The reduction in laser intensity is compensated by the increase in interaction length, as described by Eq. (4). The analytically calculated Compton yield matches well with the measured trend, as shown by the solid red curve in Fig 7(b). The measured scaling suggests that the yield compared to head-on is limited unless ultrashort laser pulses are used. However, even with a reduced interaction length, the brightness can be significantly larger due to relativistic beaming. The measured scaling also suggests control over the interaction length by laser focusing in the p direction allowing for tuning of the radiation pulse length and bandwidth.

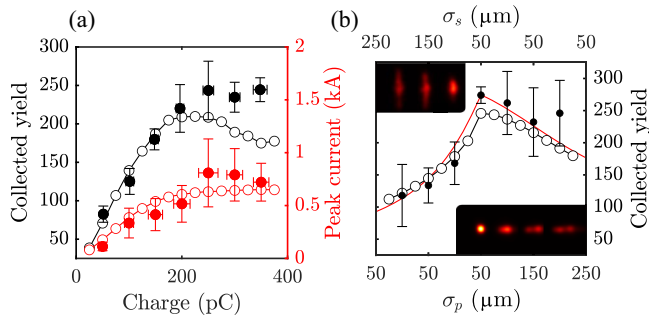


FIG. 7. (a) Collected photon number in black and peak current in red versus bunch charge as (filled) measured and (open) simulated. (b) Collected photon number versus s and p laser waist size as (filled) measured and (open) simulated. The red solid curve represents the analytical prediction given by Eq. (2), which is normalized to the data at $\sigma_{(s,p)} = 50$ μm. The insets in (b) show the transverse intensity distribution of the laser in the IP.

A. Pulse-front tilt

VI. SPATIOTEMPORAL OPTIMIZATION AND LIMITS

Although the visible photon yield produced in this specific low-energy demonstration does not constitute a practical light source, the experimental validation of the overtaking geometric scalings is invaluable for future machine design. A primary limitation of uncompensated shallow-angle Compton scattering is the restricted interaction length driven by the longitudinal transit time. In this section, we discuss a recognized technique to overcome this limitation via laser spatiotemporal shaping to address the macroscopic and nonlinear spectral broadening constraints.

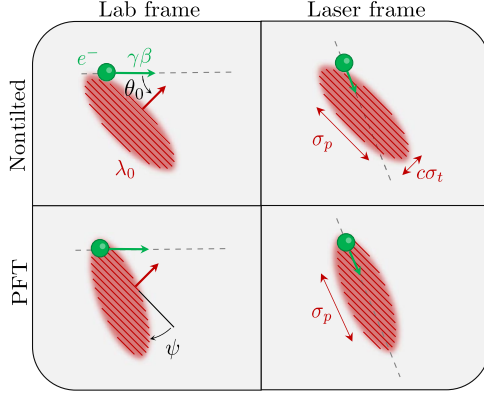


FIG. 8. Schematic of an electron interacting with an uncompensated and a pulse front tilted (velocity matched) laser pulse in the lab frame and the Galilean reference frame comoving with the laser.

The limiting contribution of the laser pulse length to the interaction time can be completely removed by imparting a pulse-front tilt (PFT) to the laser [19], a mechanism analogous to crab crossing [53]. Mathematically, PFT introduces a transverse shear to the laser density distribution in its propagation frame (p, l), effectively replacing the longitudinal envelope coordinate l with $l - p \tan \psi$ (see Fig. 8). To achieve perfect velocity matching, where the temporal center of the laser envelope co-propagates exactly with the electron bunch, the optimum tilt angle ψ is analytically found to be:

$$\tan \psi = \frac{1 - \beta \cos \theta_0}{\beta \sin \theta_0}. \quad (6)$$

Velocity-matched PFT completely removes the limiting contribution of the longitudinal transit time, extending the interaction length to $N_\varphi = k_0 \sigma_p (1 - \beta \cos \theta_0) / (\beta \sin \theta_0)$ thereby enabling a larger tunability range and enhanced intrinsic brightness [19].

To determine the exact impact of this velocity-matched PFT on the total scattered photon yield, we recalculate the luminosity integral. The out-of-plane interaction width Σ_s remains unaffected; however, the in-plane interaction width reduces to:

$$\Sigma_{p,\text{PFT}} = \left[\sigma_x^2 \left(\frac{\beta - \cos \theta_0}{1 - \beta \cos \theta_0} \right)^2 + \frac{(\sigma_z^2 + \beta^2 c^2 \sigma_t^2) \sin^2 \theta_0}{(1 - \beta \cos \theta_0)^2} \right]^{1/2}. \quad (7)$$

Comparing this to the untilted collision geometry [Eq. (4)], the velocity-matched PFT perfectly cancels the transverse laser waist σ_p contribution from the spatial overlap. Physically, because the tilted laser envelope sweeps precisely along the electron beam axis, any increase in the in-plane laser waist decreases the peak laser photon density but is precisely compensated by a

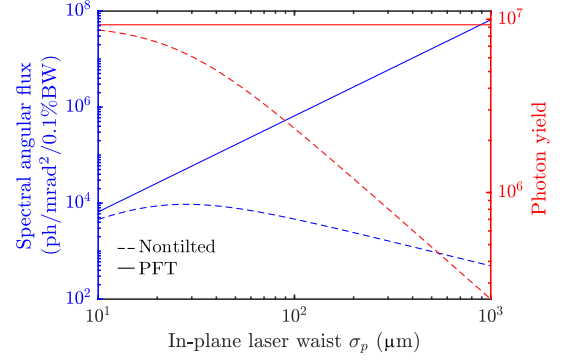


FIG. 9. The spectral angular yield (blue, left axis) and total photon yield (red, right axis) as a function of the in-plane laser waist σ_p for both a PFT laser (solid curves) and an uncompensated laser (dashed curves). The parameters correspond to a 20° collision between a 100 MeV, 100 pC, 10 fs rms electron bunch and a 1030 nm, 100 mJ, 10 fs rms laser pulse, with the electron and out-of-plane laser waists fixed at 10 μm .

proportional increase in the interaction length. This makes the total integrated yield completely independent of in-plane laser focusing. As illustrated in Fig. 9, this spatio-temporal decoupling allows the in-plane laser waist σ_p to be arbitrarily expanded to reduce the intrinsic bandwidth, driving a continuous increase in spectral angular brightness without suffering the significant drop in total yield seen in the uncompensated geometry.

During an angled collision, the finite transverse width of the electron bunch prevents all electrons from crossing the densest region of the laser pulse. The transverse beam size penalty (σ_x) therefore remains in the in-plane interaction width even with a pulse-front tilt of the laser. Crab crossing [53] can rotate the electron bunch to align its spatial profile perfectly with the tilted laser pulse front, restoring optimal spatial overlap and effectively setting the σ_x term in Eq. (5) to zero. With envelope alignment via PFT and crab crossing, the requirement of tight transverse focusing in a head-on collision is fully replaced by that of longitudinal focusing, which is extra stringent due to the small relative velocity between the beams. Consequently, while PFT successfully decouples the interaction bandwidth from the spatial laser waist, high yield still requires the use of compressed electron bunches and ultrashort laser pulses.

A. Spectral broadening limits

While velocity-matched PFT implies that the intrinsic bandwidth $\Delta\omega/\omega = 1/N_\varphi$ can be reduced indefinitely by expanding the in-plane laser waist σ_p , the achievable spectral brightness is bounded by two macroscopic constraints: inhomogeneous broadening due to the electron beam distribution and nonlinear ponderomotive forces.

First, the total radiation bandwidth is bounded by the macroscopic quality of the electron beam. The observed bandwidth along a certain angle is the quadrature sum of

the intrinsic homogeneous bandwidth and the inhomogeneous broadening driven by the electron energy spread ($\Delta\omega/\omega \simeq 2\sigma_\gamma/\gamma$) and the finite beam emittance. The effect of emittance in a side-scattering geometry is more stringent than in a head-on collision, where it scales quadratically with the beam divergence ($\Delta\omega/\omega \simeq \gamma^2\sigma_x^2$). In an overtaking geometry, the transverse angular spread σ_x directly modulates the crossing angle θ_0 , which, by differentiating the resonance condition [Eq. (1)], shows a first order term given by $(\Delta\omega/\omega)_\varepsilon \simeq \sigma_x \sin \theta_0 / (1 - \beta \cos \theta_0)$. Consequently, expanding σ_p only increases the peak spectral brightness up to an optimal threshold, defined by the point where the intrinsic bandwidth matches this inhomogeneous limit.

Second, a larger laser waist plays a critical role in mitigating nonlinear spectral broadening. To maintain a narrow, quasimonochromatic emission line, the normalized laser vector potential must be kept small, such that $a_0 \ll 1$, to prevent ponderomotive redshifting and the emission of higher-order harmonics. Scaling up the total photon yield by increasing the laser pulse energy inherently risks pushing the interaction into the nonlinear regime. In an uncompensated crossing, expanding σ_p to maintain a small a_0 could degrade the interaction length N_φ . However, with a velocity-matched PFT, the interaction length becomes independent of σ_p . This unique geometric freedom allows for the use of highly energetic laser pulses distributed over a massive in-plane waist, maximizing the total yield while simultaneously suppressing nonlinear broadening and optimizing the spectral brightness.

Beyond the limits imposed by inhomogeneous broadening, implementing velocity-matched PFT introduces specific geometric and optical constraints that can influence the radiation spectrum. Generating PFT requires the use of reflective gratings to deliberately introduce angular dispersion. However, because propagating an angularly dispersed beam inherently causes the pulse to stretch in time, the laser must be prechirped such that it reaches optimal longitudinal compression, i.e., becomes fully group delay dispersion free, exactly at the interaction plane [54]. Furthermore, achieving high photon density requires focusing this spatially dispersed pulse along the electron trajectory. The coupling of the angular dispersion and the focusing optics can lead to higher-order wave-front aberrations that need to be carefully managed. Finally, these setups demand rigorous alignment and stability to prevent spectral variations in the laser pulse. A framework for designing these setups and a detailed discussion on the practical limitations are provided in Ref. [54].

VII. PROSPECTS FOR HIGH-BRIGHTNESS ARCHITECTURES

To highlight the geometric scaling behavior of shallow-angle ICS, we consider three design cases targeting

distinct x-ray and gamma-ray regimes, driven by different accelerator technologies: a conventional radio-frequency linear accelerator (linac) [20], a laser wakefield accelerator (LWFA) [55], and a diffraction-limited storage ring (synchrotron) [56]. The parameters are summarized in Table I. For all cases, we assume a collision driven by a laser equipped with velocity-matched PFT, providing an average power of 1 kW to ensure a balanced comparison across the different architectures. The expressions of the brightness and source sizes used for the calculations are given in Appendix A.

The Linac case represents a mature, high-stability architecture capable of delivering low-emittance, 100 MeV electron bunches [20]. In a traditional head-on collision with a 1030 nm laser, generating x-rays with a photon energy around 12 keV would require a significantly lower electron beam energy of 26 MeV, drastically expanding the radiation cone by the weaker relativistic beaming. By utilizing a 30° shallow-angle geometry, we can produce the same x-rays with the 100 MeV beam energy design yielding a highly collimated x-ray beam with a central opening angle of just 0.328 mrad. While a symmetric head-on collision at the lower energy would produce a total photon yield approximately an order of magnitude higher, its spectral brightness is more than an order of magnitude lower. Due to the high 100 kHz repetition rate, this configuration delivers an average spectral brightness of 6.1×10^{13} ph/(s mm² mrad² 0.1% BW). This provides a highly collimated x-ray source comparable to synchrotron bending magnets, but realized within a highly accessible, out-of-line collision geometry.

The LWFA case demonstrates why this accelerator architecture is arguably the perfect driver for shallow-angle ICS. Operating at 100 MeV represents a highly favorable operating condition for modern LWFAs, offering robust stability, high beam quality, and reliable charge extraction from a plasma target from modest laser drivers [55]. The shallow-angle geometry allows us to utilize this optimal electron energy to directly generate 15.9 keV x-rays, while simultaneously solving a fundamental limitation of head-on collisions driven by high-peak-power lasers. The LWFA utilizes a 1 J, 5 fs Ti:Sapphire laser. If this extreme pulse were focused tightly in a head-on geometry, the normalized vector potential (a_0) would greatly exceed unity, driving the interaction deep into the nonlinear regime and severely broadening the emission spectrum. However, in the shallow-angle setup, we are able to expand the waist significantly in the crossing plane $\sigma_p = 2000$ μm lowering the intensity to the linear regime, preserving the narrow bandwidth while utilizing the full 1 J of pulse energy. Furthermore, because the LWFA naturally delivers an ultrashort 2 fs electron bunch, the longitudinal penalty that normally plagues shallow-angle yield is essentially nonexistent. The intense laser and electron bunch overlap densely, driving a high peak brightness

TABLE I. Design parameters for shallow-angle ICS sources across three distinct electron-driver architectures, assuming velocity-matched PFT and a 1 kW average power laser system.

Parameter	Symbol	Design case		
		Linac	LWFA	Synchrotron
Electron beam				
Beam energy (MeV)	E_e	100	100	2000
Charge (pC)	Q	250	100	1000
Repetition rate (Hz)	f_{rep}	1.0×10^5	1.0×10^3	1.0×10^6
Bunch length (rms) (ps)	σ_z/c	0.05	0.002	110
Energy spread (rms)	σ_γ/γ	0.1%	2%	0.098%
Normalized emittance ($\mu\text{m}\cdot\text{rad}$)	ϵ_n	1	1	0.27
Waist at IP (rms) (μm)	$\sigma_{x,y}$	10	5	10
Compton laser (with PFT)				
Wavelength (nm)	λ_0	1030	800	800
Pulse duration (rms) (fs)	σ_t	50	5	50
Pulse energy (mJ)	E_c	10	1000	1
Average power (kW)	P_c	1.0	1.0	1.0
Crossing angle ($^\circ$)	θ_0	30	30	11.8
Out-of-plane waist (rms) (μm)	σ_s	10	5	10
In-plane waist (rms) (μm)	σ_p	200	2000	200
Normalized vector potential	a_0	0.022	0.244	0.005
Radiation properties				
Photon energy (keV)	E_{ph}	12.4	15.9	1003
Total yield [ph/shot]	N_{ph}	7.6×10^5	4.8×10^8	5.9×10^1
Average flux (ph/s)	Φ_{avg}	7.6×10^{10}	4.8×10^{11}	5.9×10^7
Pulse duration (rms) (fs)	$\sigma_{t,\text{rad}}$	49.5	2.0	145.8
Source size x/y (rms) (μm)	$\sigma_{x,y,\text{rad}}$	10.0/7.1	5.0/3.5	10.0/7.1
Total bandwidth (rms)	$(\Delta\omega/\omega)_{\text{tot}}$	0.41%	4.02%	0.65%
Opening angle (central) (mrad)	$\sigma_{\theta,\text{rad}}$	0.328	1.02	0.021
Opening angle (full BW) (mrad)	$1/\gamma$	5.11	5.11	0.26
Peak brightness [ph/(s mm ² mrad ² 0.1% BW)]	B_{peak}	4.9×10^{21}	3.3×10^{24}	2.1×10^{19}
Average brightness [ph/(s mm ² mrad ² 0.1% BW)]	B_{avg}	6.1×10^{13}	1.6×10^{13}	7.7×10^{12}

of 3.3×10^{24} ph/(s mm² mrad² 0.1% BW), rivaling modern undulators from an ultracompact footprint.

Finally, the Synchrotron case illustrates how shallow-angle ICS can extend the capabilities of a 2.0 GeV multibend achromat storage ring into an otherwise inaccessible spectral (and temporal) regime. Colliding the beam with a highly repetitive Ti:Sapphire laser (800 nm, 1 mJ, 50 fs) operating at 1 MHz at an 11.8° crossing angle generates 1 MeV γ -rays, a spectral range that is inaccessible to undulator radiation at this beam energy. Furthermore, the electron bunch length (110 ps) is vastly longer than the 50 fs laser pulse. In this extreme disparity regime, the laser pulse effectively slices through a microscopic fraction of the electron bunch. While the total yield per pass is severely restricted by the spatial overlap, the geometry generates an ultrashort 150 fs γ -ray pulse from an otherwise macroscopic storage ring bunch, highlighting the utility of shallow-angle ICS as a high-energy femtosecond slicing tool [23]. Moreover, the radiation pulse

length can be tuned (shorter or longer) by the waist size of the pulse front tilted laser. The increased collimation results in a decent brightness at the high photon energy. Operating at a high repetition rate with a low laser pulse energy resulting in a small number of photons per shot is useful for applications such as nuclear resonance fluorescence [57,58], which rely on precise photon counting limited by detector pileup [59].

VIII. CONCLUSION

We have presented a comprehensive theoretical, experimental, and numerical investigation of shallow-angle inverse Compton scattering. While overtaking collision geometries introduce specific spatiotemporal challenges, such as stringent sensitivity to time-of-arrival jitter and a longitudinal luminosity penalty, our experimental results at the UCLA Pegasus photoinjector confirm that these dynamics are well understood and predictable. Specifically, the

measured radiation properties validate our analytical scaling laws, including the asymmetric dependence of the photon yield on the transverse laser dimensions and the yield saturation effects driven by decompression of the electron bunch due to space charge.

Having benchmarked the fundamental interaction, we show how pulse-front tilt can circumvent interaction length limits of shallow-angle Compton, which can improve the spectral brightness significantly. Furthermore, this spatiotemporal shaping enables optimal use of high powered laser pulses while remaining in the linear regime, circumventing the nonlinear ponderomotive broadening that limits head-on collisions. Including this to the analytical framework and applying it to future light source designs, we showed that the shallow-angle regime is highly synergistic with modern accelerator technologies. Collectively, these findings establish the shallow-angle geometry as a robust, scalable, and highly tunable pathway for the next generation of compact, high-brightness light sources.

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DATA AVAILABILITY

The data that support the findings of this article are not publicly available. The data are available from the authors upon reasonable request.

APPENDIX A: RADIATION BRIGHTNESS AND SOURCE DIMENSIONS

The peak and average spectral brightness of the scattered radiation are critical figures of merit for evaluating source performance, defined, respectively, as

$$B_{\text{peak}} = \frac{N_{\text{ph}}}{(2\pi)^{5/2} \sigma_{x,\text{rad}} \sigma_{y,\text{rad}} \sigma_{\theta,\text{rad}}^2 \sigma_{t,\text{rad}}} \frac{0.1\%}{(\Delta\omega/\omega)_{\text{tot}}},$$

$$B_{\text{avg}} = B_{\text{peak}} \sqrt{2\pi} \sigma_{t,\text{rad}} f_{\text{rep}},$$

where N_{ph} is the total photon yield per shot, f_{rep} is the interaction repetition rate, and $(\Delta\omega/\omega)_{\text{tot}}$ is the total relative bandwidth (rms). The central opening angle (rms) of the radiation is given by $\sigma_{\theta,\text{rad}} = \sqrt{(\Delta\omega/\omega)_{\text{tot}}}/\gamma$. The temporal duration of the emitted x-ray pulse $\sigma_{t,\text{rad}}$, as well as the transverse spatial source dimensions $\sigma_{x,\text{rad}}$ and $\sigma_{y,\text{rad}}$, are governed strictly by the spatiotemporal convolution of the electron and laser distribution envelopes during the collision.

The spatiotemporal volume of the radiation generation is defined by the product of the electron and laser Gaussian density distributions: $\rho_{\text{rad}} = \rho_e \rho_l$. By calculating the moments of the radiation distribution $\langle X^n \rangle = \int \rho_x X^n dt d^3r$, we can analytically determine the physical source size $\sigma_{X,\text{rad}} = (\langle X^2 \rangle - \langle X \rangle^2)^{1/2}$ of the scattered x-ray pulse. In the following we do not consider pulse-front tilt of the laser.

Because the out-of-plane ($X = y$) source size is completely decoupled from the crossing dynamics, its standard deviation is simply the harmonic mean of the initial transverse beam sizes given by $\sigma_{y,\text{rad}} = (1/\sigma_y^2 + 1/\sigma_s^2)^{-1/2}$. The in-plane dimensions ($X = x$) are coupled by the scattering angle θ_0 and the relativistic velocity of the electrons. The in-plane source size is given by

$$\sigma_{x,\text{rad}} = \sigma_x \left[1 - \frac{\sigma_x^2}{\Sigma_p^2} \left(\frac{\beta - \cos \theta_0}{1 - \beta \cos \theta_0} \right)^2 \right]^{1/2}, \quad (\text{A1})$$

where Σ_p is the effective in-plane dilution dimension defined in Eq. (3). These expressions show that the source size is determined by the combination of the individual beam widths and lengths, bounded by their respective size. The duration of the emitted flash ($X = t - z/c$) as seen by a distant observer in the laboratory frame, evaluates to:

$$\sigma_{t,\text{rad}} = \frac{1}{c \Sigma_p (1 - \beta \cos \theta_0)} \times [\sigma_z^2 ((\sigma_x^2 + \sigma_p^2)(1 - \cos \theta_0)^2 + c^2 \sigma_t^2 \sin^2 \theta_0) + (1 - \beta)^2 (c^2 \sigma_t^2 \sigma_p^2 + \sigma_x^2 \cos^2 \theta_0) + \sigma_x^2 \sigma_p^2 \sin^2 \theta_0]^{1/2}. \quad (\text{A2})$$

For the PFT case the source size in y remains the same. In x the source size can be found by making the substitution $\Sigma_p \rightarrow \Sigma_{p,\text{PFT}}$ in Eq. (A1). The pulse length in this case is given by

$$\sigma_{t,\text{rad,PFT}}^2 = \frac{1}{c^2 \Sigma_{p,\text{PFT}}^2 (1 - \beta \cos \theta_0)^2} \times [\sigma_z^2 (\sigma_x^2 (1 - \cos \theta_0)^2 + c^2 \sigma_t^2 \sin^2 \theta_0) + (1 - \beta)^2 c^2 \sigma_t^2 \sigma_x^2 \cos^2 \theta_0] + \left(\frac{(1 - \beta) \sigma_p}{\beta c \sin \theta_0} \right)^2. \quad (\text{A3})$$

APPENDIX B: WIENER DECONVOLUTION

The Compton time traces are affected by random noise from background radiation and shifts in time due to TOA jitter between the laser and electron beam. The observed signal can be written as $y(t) = f(t) * h(t) + \eta(t)$, where t here is the delay imparted by the IR delay line, $f(t)$ is the true Compton scattering time trace, $h(t)$ is the TOA

jitter function, and $\eta(t)$ describes the random background noise. Here we assume that the jitter is Gaussian such that $h(t) = \exp[-t^2/(2\sigma_{\text{toa}}^2)](2\pi\sigma_{\text{toa}}^2)^{-1/2}$ with σ_{toa} the time of arrival jitter.

In Wiener deconvolution [60], the Fourier transform of the original signal $F(\omega) = \mathcal{F}\{f(t)\}$, where $\mathcal{F}$ denotes the Fourier transform, is approximated by $F(\omega) = W(\omega)Y(\omega)$, where $Y(\omega) = \mathcal{F}\{y(t)\}$ is the Fourier transform of the observed signal and $W(\omega)$ is the Wiener filter given by $W(\omega) = H^*(\omega)/[|H(\omega)|^2 + 1/\text{SNR}(\omega)]$, where $H(\omega) = \mathcal{F}\{h(t)\}$ the Fourier transform of the jitter function and SNR the (estimated) signal-to-noise ratio. For vanishing noise ($\text{SNR} \rightarrow \infty$), the Wiener filter is simply the inverse of the spectral jitter function, which deconvolves the signal in the expected way. For nonzero noise the second term in the denominator attenuates frequencies according to the signal-to-noise ratio.

To find the time of arrival jitter we apply a Wiener filter to the low noise EOS time trace, for varying rms time of arrival jitter σ_{toa} . The SNR used in the Wiener filter was estimated in this case from the ratio between the integrated power spectrum below (signal) and above 4 THz (noise). We find that the average spread of the observed EOS signal was minimized for an rms time of arrival jitter of 400 fs, which is consistent with the rf fluctuations in phase and amplitude.

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