

High-order harmonic generation from ultrarelativistic s-polarized 400 nm laser-solid interactions

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Harmonic emission from s-polarized, 400 nm, ultrahigh-contrast laser-solid interactions is observed to exhibit a plasma-frequency cutoff at relativistic intensities ($\sim 10^{20}$ W cm⁻²), with harmonics emerging beyond this cutoff in the ultrarelativistic regime ($\sim 10^{21}$ W cm⁻²). Experimental conditions isolate $\vec{J} \times \vec{B}$ heating as the dominant absorption mechanism, and particle-in-cell simulations show that its intensity-dependent evolution drives distinct harmonic generation mechanisms.

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The interaction of high-intensity lasers with solid-density plasmas represents a fundamental nonlinear regime of light-matter coupling, with applications ranging from compact ion acceleration and directed neutron generation to x-ray radiography [1–4]. These interactions also drive high-order harmonic generation (HHG) [5], producing coherent soft- to hard-x-ray radiation with attosecond (1 as = 10^{-18} s) pulse durations [6]. Notably, solid-density plasmas impose no practical upper limit on the applied laser intensity [7], enabling direct scaling to higher photon energies [8,9]. Control over attosecond emission has been demonstrated using two-color fields [10], curved relativistic mirrors [11], and single-pulse isolation schemes [12–14]. Combined with high conversion efficiencies [15,16] and strong spatial coherence [15], solid-density HHG provides a versatile source of ultrafast radiation for studies in atomic [17,18] and biomedical systems [19].

The emission of coherent harmonic radiation arises from a nonlinear reflection at the overdense plasma-vacuum interface in which laser light is up-converted to integer multiples of the fundamental frequency. Theoretical and experimental studies demonstrate two leading mechanisms for HHG: coherent wake emission (CWE) [20] and the relativistic oscillating mirror (ROM) [21–23]. In the CWE process, electron bunches are accelerated into the interface, exciting plasma oscillations in their wake every optical cycle. These oscillations produce

coherent radiation via linear mode conversion up to the plasma frequency, $\omega_p = \sqrt{n_e e^2 / m_e \epsilon_0}$, where n_e , e , m_e , and ϵ_0 denote the peak electron density, electron charge, electron mass, and vacuum permittivity, respectively. Because this process is dominated by plasma oscillations, the harmonic spectrum depends only weakly on the laser intensity, with experiments demonstrating harmonic generation at laser intensities as low as 6.25×10^{15} W cm⁻² with an 800 nm laser beam [20]. The maximum harmonic order scales as $N_{\max} = \sqrt{n_e / n_c}$, where $n_c = \epsilon_0 m_e \omega_L^2 / e^2$ is the critical density for a laser of frequency ω_L . Conversely, the ROM mechanism is strongly intensity dependent and arises only when electrons attain relativistic motion. This becomes dominant when the normalized vector potential $a_0 = eE_0 / m_e \omega_L c \geq 1$, where E_0 is the laser electric field and c is the speed of light in vacuum. The relativistic critical surface induces periodic phase modulations in the reflected laser field, generating harmonic radiation beyond the plasma frequency [23].

Experimental observations of HHG require ultrahigh temporal laser contrast, defined as the ratio of prepulse to main laser pulse intensity, to mitigate long-scale preplasma formation, which can suppress harmonic generation [15,16]. Contrast can be improved through means such as cross-polarized wave generation [24] or relativistic plasma shutters [7,25–32]. More recently, significant contrast improvements have been achieved by frequency doubling the fundamental laser output, which, unlike plasma shutters, has the practical advantage of enabling operation at the full repetition rate of the laser system [16,33–37]. This technique has now been extended to petawatt-class lasers [38]. In parallel, most HHG studies have employed low-Z dielectric targets such as fused silica and plastics, which are generally less sensitive to laser contrast than metals [39]. These materials are also more readily driven to fully ionized states at moderate intensities, enabling a clear

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identification of the spectral cutoff associated with the plasma frequency of the fully ionized target [7].

Further mitigation of the preplasma thermal expansion may be facilitated by reducing electron temperatures through the suppression of collisionless absorption mechanisms. This can be accomplished by employing oblique incidence with *s*-rather than *p*-polarization [40–42], which removes the laser field component normal to the target surface, thereby suppressing both resonant [43] and Brunel/vacuum [44] heating. For *s*-polarized interactions, laser absorption is dominated by $\vec{J} \times \vec{B}$ heating [45], which becomes efficient only in the relativistic regime. As a result, observations of robust harmonic generation in *s*-polarized interactions have been limited [46], leaving its role in HHG largely unexplored experimentally. Consequently, most studies have been conducted using *p*-polarized interactions, exploiting the additional collisionless absorption mechanisms to efficiently drive HHG and revealing the transition from a resonant-absorption-like mechanism to one of relativistic origin [37].

In this Letter, we report high-order harmonic radiation in the specular reflection of an ultrahigh-contrast, *s*-polarized, 400 nm laser interacting with solid targets at relativistic intensities. Two distinct regimes are identified. In the relativistic regime ($a_0 \sim 6$), harmonics extend to a well-defined cutoff corresponding to the plasma frequency of a nearly fully ionized copper target, a spectral feature not previously resolved in high-*Z* materials, while in the ultrarelativistic regime ($a_0 \sim 15$), the spectrum extends beyond. Particle-in-cell (PIC) simulations demonstrate that $\vec{J} \times \vec{B}$ heating drives two intensity-dependent HHG mechanisms through strong laser-plasma coupling at the relativistic critical-density surface.

The experiment was conducted using the 400 nm beamline at the Advanced Laser for Extreme Photonics, a petawatt-class Ti:sapphire laser facility. The beam was produced by frequency doubling the 800 nm, 185 mm diameter, near-flat-top fundamental in a 0.8-mm-thick antireflection-coated KH_2PO_4 crystal. The residual fundamental was filtered using a sequence of five dichroic mirrors, yielding an *s*-polarized, ultrahigh-contrast pulse [38]. *P*-polarized interactions were not explored due to size constraints of the polarization-rotating optics. The mean temporal contrast of the fundamental was measured from a third-order autocorrelation trace to be $\sim 10^7$ on the nanosecond timescale preceding the main pulse. This corresponds to an inferred contrast of $> 10^{13}$ for the resulting 400 nm beam, spanning conditions from below to near frequency-doubling saturation. The laser pulse, with a duration of 45 fs full width at half maximum (FWHM), was focused using an $f/2$ off-axis parabolic (OAP) mirror to a Gaussian-like spot with a diameter of $1.60 \pm 0.02 \mu\text{m}$ FWHM (Fig. 1). Under this configuration, laser energies of 1.13 ± 0.01 and 6.22 ± 0.11 J were delivered to the target for high-intensity (HI) and ultrahigh-intensity (UHI) shots, corresponding to mean normalized vector potentials $a_0 = 6.47 \pm 0.04$ and $a_0 = 15.17 \pm 0.15$, respectively.

The primary targets were freestanding copper (Cu) foils of thicknesses 1, 5, and $10 \mu\text{m}$, oriented at a 22.5° angle of incidence relative to the laser axis. Glass and layered Cu-plastic targets were also used, but are not shown here. The foils were translated to provide a fresh target for each shot

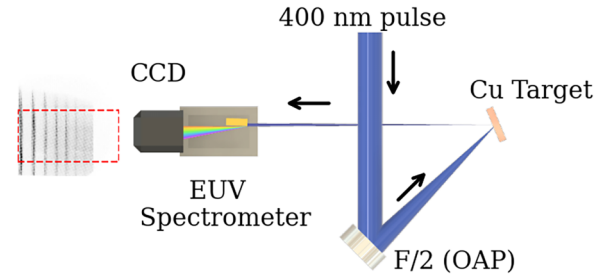


FIG. 1. Experimental setup for measuring radiation in the specular reflection from Cu targets. A representative charge-coupled device (CCD) image shows the radiation lines; dashed rectangle marks the analysis region.

and monitored with an objective camera to ensure alignment with the laser focal plane. With controlled target positioning, additional out-of-focus shots at 6.08 ± 0.05 J were taken. The non-Gaussian focal spot was characterized using an encircled-energy method, with 50% of the total signal enclosed within a diameter of $8.8 \mu\text{m}$, corresponding to a mean $a_0 = 3.55 \pm 0.01$. In total, nine shots were taken per foil thickness: ~ 1 and ~ 6 J delivered at best focus, and ~ 6 J delivered out of focus.

Several diagnostics, including a von Hamos mica crystal spectrometer and a high-resolution spectrometer consisting of two spherical germanium crystals, were used to measure copper K-shell emission, while a magnetic electron spectrometer measured rear-side electron energies. The results obtained with the crystal spectrometers are reported in Ref. [47]. The home-built extreme ultraviolet spectrometer [48] was aligned to the specular reflection defined by the plane of incidence, where it remained fixed; consequently, any azimuthal rotation of the target to modify the incident polarization would redirect the specular beam away from any available diagnostic ports. The detector captured wavelengths from 14 to 25 nm, where the spectral resolution was limited by the $13 \mu\text{m}$ pixel size of the charge-coupled device. The effective wavelength range was restricted to 17–25 nm by a set of aluminum x-ray filters, which blocked the reflected laser light. The first Al filter, mounted on a 15-mm-diameter aperture, was positioned ~ 2 m from the source at the spectrometer entrance, corresponding to an acceptance solid angle of 3.91×10^{-5} sr after accounting for partial clipping.

Radiation at both even and odd multiples of the fundamental frequency was consistently observed from Cu foils irradiated at HI and UHI, whereas the defocused shots frequently produced no harmonics. Glass targets did not produce harmonic emission within the detector range. Alongside integer-harmonic peaks, a continuum was present in all shots (Fig. 1, left). Figures 2(a) and 2(b) show representative response-corrected spectra measured at HI and UHI, respectively, for a $5\text{-}\mu\text{m}$ target, plotted as photons detected versus harmonic order.

At HI, the harmonic spectrum extends up to the 17th order, exhibiting a pronounced spectral cutoff. Before this cutoff, harmonic radiation follows a mean power-law scaling of $n^{-11.98}$. Beyond the cutoff, more than half of the shots exhibited the emergence of radiation near the 18th harmonic, with flux barely distinguishable above the continuum [Figs. 2(a) and 2(c)]. The mean peak photon count for the 17th harmonic

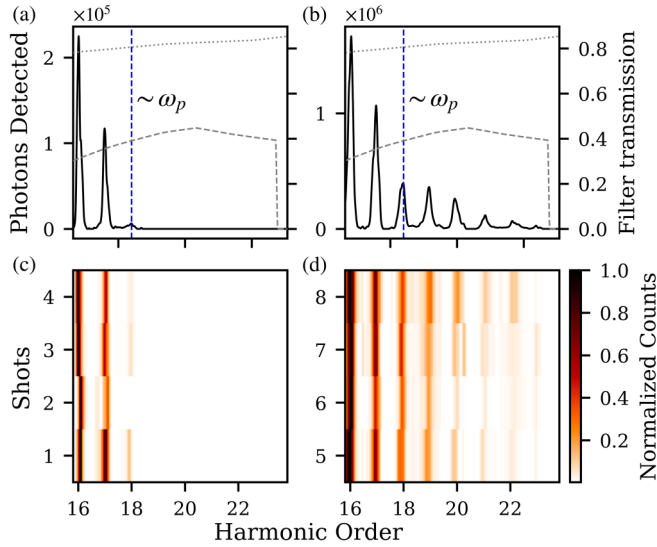


FIG. 2. Response-corrected lineouts of harmonic spectra from 5- μm -thick Cu foils at $a_0 = 6.4$ [panel (a)] and 15.1 [panel (b)]. Dashed and dotted curves show the transmission of 0.6- μm Al filter and CCD quantum efficiency, respectively. The dashed blue line marks the estimated maximum plasma frequency. (c), (d) Waterfall plots of harmonic spectra from Cu targets of varying thickness at HI (shots 1–4: 10, 10, 1, and 10 μm) and UHI (shots 5–8: 5, 1, 1, and 1 μm).

(23.5 nm) was $8.87 \pm 4.44 \times 10^4$ photons per shot, with substantial shot-to-shot variation and a maximum of 1.57×10^5 .

At UHI, harmonic radiation is observed near the aluminum edge at 17.2 nm, with a maximum harmonic reaching the 23rd order [Figs. 2(b) and 2(d)]. Harmonics above the 18th order follow a separate mean power-law scaling of $n^{-9.82}$. The mean peak photon count for the 17th harmonic was $5.26 \pm 2.89 \times 10^5$ photons per shot, approximately six times higher than that in the HI case, with a maximum of 1.07×10^6 . As shown in the waterfall plot [Fig. 2(d)], the absolute signal decreases as the harmonic order increases.

For the defocused shots, cases exhibiting harmonic radiation show a consistent cutoff at the 17th harmonic order, though not shown here. The mean peak photon count for the 17th harmonic is reduced to $7.27 \pm 2.16 \times 10^4$ photons per shot, comparable to the HI shots. The elevated continuum background and limited number of recorded shots preclude a reliable determination of their scaling.

Given the combination of relativistic intensities ($I \geq 8 \times 10^{18} \text{ W cm}^{-2}$ for 400 nm wavelength) and s -polarized, shallow-incidence interactions, resonance absorption and vacuum heating can be assumed to be suppressed, leaving relativistic $\vec{J} \times \vec{B}$ as the dominant mechanism for electron heating. Collisionless absorption processes are critical for laser-plasma coupling and thus directly impact harmonic generation mechanisms [23,49]. For p -polarized interactions, the electric field component perpendicular to the target directly excites several absorption processes, enabling harmonic generation below the relativistic threshold, as in the CWE mechanism, while enhancing laser energy absorption into the plasma [41,42] and potentially reducing the power into the

harmonics [50]. In contrast, s -polarized interactions are dominated by $\vec{J} \times \vec{B}$ heating and are therefore often considered unfavorable for HHG at moderate intensities, where harmonics exhibit low conversion efficiency and few, if any, harmonic orders are produced [15,16,37,40].

Under $\vec{J} \times \vec{B}$ heating, the longitudinal component of the magnetic term in the Lorentz force drives plasma oscillations at $2\omega_L$ and scales as a_0^2 [51]. Consequently, sufficiently high laser intensities may be required to efficiently drive plasma waves throughout the density gradient via energetic electron bunches, analogous to the CWE mechanism. The crucial difference is that the process becomes intensity dependent, with CWE-like harmonics potentially requiring laser strengths of $a_0 \geq 1$. A clear signature of this regime is a sharp cutoff at the maximum plasma frequency. At still higher, ultrarelativistic intensities, strong surface oscillations may dominate the interaction and modulate the reflected pulse, producing ROM harmonics. In that case, harmonic orders extending beyond the plasma frequency would be consistent with the onset of ROM. In addition to spectral differences, previous experiments have distinguished CWE and ROM harmonics through their angular emission profiles [52]. ROM harmonics are highly collimated about the reflection axis, approaching a diffraction-limited divergence at relativistic intensities [32], with harmonic divergence scaling as $\theta_n \sim \theta_{\text{laser}}/n$, where θ_{laser} is the divergence of the laser beam and n is the harmonic order. In contrast, CWE harmonics exhibit larger divergence due to their intrinsic phase [53]. Assuming diffraction-limited emission, the highest harmonic order ($n = 23$) has an estimated solid angle of $\Omega_{23} \approx \pi(\theta_{23})^2 \approx 5.93 \times 10^{-5} \text{ sr}$, exceeding the acceptance of the spectrometer; consequently, the harmonic angular profiles could not be resolved. Instead, the mechanisms are distinguished by the plasma-frequency cutoff and the differing decay rates.

With these signatures in mind, copper ionization data from NIST [54], together with the barrier-suppression ionization (BSI) model [55,56], yield a maximum charge state of Cu^{27+} (fully ionized L shell) at HI and UHI, corresponding to a peak electron density of $\sim 329n_c$ and a CWE cutoff near the 18th harmonic. This idealized estimate, based on strictly field-driven ionization, neglects additional processes such as collisional ionization [57] and the strong spatial and temporal gradients inherent to solid-density plasma interactions. Nevertheless, the estimated high degree of ionization and corresponding plasma density are consistent with values inferred from x-ray crystal spectrometer measurements on the front surface of copper foils of comparable thickness. These measurements report ionic line emission at $\sim 10^{21} \text{ W cm}^{-2}$, together with $K\alpha$ and $K\beta$ emission at $\sim 10^{20} \text{ W cm}^{-2}$, indicative of a hot, highly ionized surface plasma [47]. Experimentally, the spectra exhibit a consistent cutoff at the 17th order when irradiated at HI, corresponding to a BSI-limited charge state of $\text{Cu}^{25\pm 1+}$ and a CWE cutoff of 17.6 ± 0.5 . The slight reduction from the predicted value can be attributed to focal averaging and spatial intensity variations, which reduce the effective plasma density. Accordingly, the observed cutoff is consistent with a plasma-frequency-limited CWE spectrum from a nearly fully ionized Cu plasma, with absorption governed by relativistic $\vec{J} \times \vec{B}$ heating rather than conventional vacuum heating. To our knowledge, such a reproducible cutoff

from high- Z materials has not been previously reported [34], primarily due to incomplete ionization states associated with their higher atomic number [7,30]. Therefore, the combination of sufficiently high laser intensities, enabling near-complete copper ionization, and ultrahigh laser contrast, which ensures a relatively short plasma scale length, is essential for resolving the plasma-frequency-limited cutoff.

Although not directly measured, short preplasma scale lengths are supported by previous experiments conducted at the same laser facility using nanostructure targets under comparable laser conditions [58,59], where preservation of the surface structure requires minimal preplasma formation. Direct control of the preplasma scale length has been demonstrated for p -polarized interactions, enabling optimal harmonic generation efficiencies either through systematic prepulse scans [52,60] or via fine-tuning of the temporal profile on the subpicosecond scale [61]. Such optimization was not available in the present experiment and may contribute to the absence of detectable harmonics from glass targets within the detector range. The ultrahigh laser contrast, combined with the higher laser fluence threshold of dielectrics, may lead to a substantially steep plasma density gradient and consequently reduce the efficiency of ROM harmonics [52,60,62], the only harmonic orders accessible to the detector for glass targets, given that the CWE cutoff corresponds to the ~ 10 th harmonic order. A systematic investigation of preplasma scale length effects on $\vec{J} \times \vec{B}$ -driven harmonics under s -polarization remains of interest but lies beyond the scope of the present work.

Harmonic orders are observed to extend beyond the cutoff only at the UHI ($a_0 \sim 15$) regime, suggesting the onset of a different mechanism, namely, ROM. This spectral transition was observed to be independent of target thickness [Figs. 2(c) and 2(d)]. Additionally, the distinctly shallower mean decay rate of harmonics above the 18th order is qualitatively consistent with the expectation that ROM operates more efficiently in the relativistic regime. The power scalings characterize the spectral decay below and above the expected plasma-frequency cutoff, rather than indicating a universal scaling as predicted by the gamma spike theory [63], where the ROM harmonic spectrum follows an $n^{-8/3}$ dependence up to a critical frequency that scales as $\gamma_{\max}^3$, beyond which it decays exponentially. Here, the combination of high plasma density and relatively short plasma scale length leads to a strong restoring force provided by the background ions, which significantly limits the maximum gamma factor of the surface electrons, placing this roll-off toward lower ROM orders [60]. Therefore, the measured spectra, which only capture the first five ROM copper harmonics, may already reside in the exponential decay section, leading to the faster decay rates measured.

To elucidate the intensity dependence of CWE and ROM harmonics in s -polarized interactions, and thereby isolate the role of $\vec{J} \times \vec{B}$ heating in both mechanisms, we performed fully relativistic PIC simulations with the code Smilei [64]. The oblique interaction was modeled using a one-dimensional (1D) boosted-frame approach [22,65–67]. The simulations were performed on a grid with high spatial resolution, $\lambda/512$, and with 180 particles per cell for each species. A laser pulse entered from the vacuum side and interacted with a fully

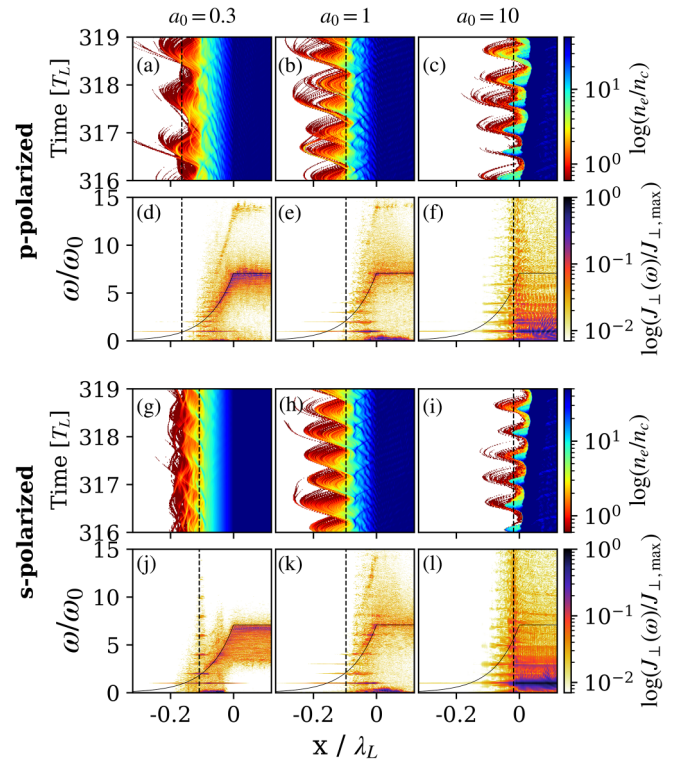


FIG. 3. PIC simulations of p -polarized [panels (a)–(f)] and s -polarized [panels (g)–(l)] interactions showing the temporal evolution of plasma density and Fourier-transformed transverse currents for a $50n_c$ target at $a_0 = 0.3, 1$, and 10 . Dashed black lines indicate the position of the relativistic critical surface.

ionized Cu plasma ($m_i/Zm_e \sim 4055$), featuring an exponentially increasing preplasma density profile, $n_e = n_{e0}\exp(x/L)$ for $x < 0$, followed by $1\text{-}\mu\text{m}$ homogeneous region ($x > 0$), where n_{e0} is the initial solid density and L is the scale length. Corresponding simulations for p -polarized interactions were also performed for comparison.

The simulations employed a reduced plasma density of $50n_c$ and a tenth-order super-Gaussian laser profile to minimize envelope effects, enabling well-resolved plasma-surface behavior while retaining the essential physics. The preplasma scale length was kept fixed at $L = \lambda/20$, with the normalized laser strength, a_0 , being the only parameter varied. Sufficient vacuum space was available to capture the entire reflected pulse for constructing the time-integrated spectra. At oblique incidence, selection rules predict both odd and even harmonics for s - and p -polarized interactions [22]. In the p -polarized case, the emission remains p -polarized [Figs. 4(a)–4(c), blue]. In contrast, for s -polarized interactions, the odd harmonics are s -polarized and are associated with the J_z current [Figs. 4(d)–4(f), red (TE)], while the even harmonics are p -polarized and are associated with the J_y current [Figs. 4(d)–4(f), black (TM)] [22,23]. The temporal evolution of the electron density was tracked over the entire simulation as a function of position normalized to the laser wavelength. A selected region at the plasma-vacuum interface, spanning four laser periods, captured the time window in which all harmonics are present. The transverse electron current densities, responsible for the emission of harmonic radiation, were

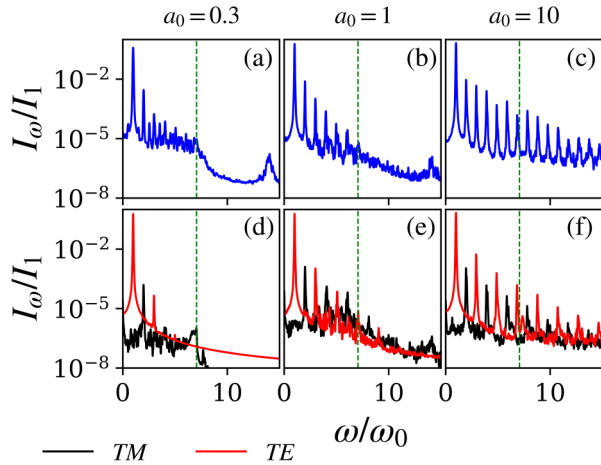


FIG. 4. Time-integrated spectra for p -polarized [panels (a)–(c)] and s -polarized [panels (d)–(f)] interactions corresponding to the PIC simulations in Fig. 3. Green dashed line indicates the CWE cutoff at ω_p .

Fourier transformed in time, $J_{\perp}(\omega)$, and plotted as a function of normalized position within the same spatial range.

At moderate intensities ($a_0 = 0.3$), the $\vec{J} \times \vec{B}$ force is resonant at $\omega_p = 2\omega_L$, enabling the laser pulse to propagate up to a critical density of $n_{c(2\omega_L)} = 4n_c$ before reflection. Along this critical surface, dense electron oscillations develop within the density gradient, occurring twice per optical cycle [Fig. 3(g)]. Fourier-transformed transverse currents reveal electron bunches centered around this region, forming a current plane whose amplitude decreases with increasing frequency [Fig. 3(j)]. Only the currents corresponding to the first three harmonic orders extend inside the density gradient. A harmonic cutoff near the fourth order, well below the CWE-predicted $N_{\max} = 7$, is observed [Fig. 4(d)]. In contrast, for p -polarization, the laser electric field drives large-amplitude density jets that oscillate once per optical cycle and propagate into the bulk plasma [Fig. 3(a)]. The corresponding current bunches are excited throughout the density gradient and a pronounced structure at $2\omega_p$ ($\sim 14\omega_L$) is observed [Fig. 3(d)]. Accordingly, the spectrum extends up to the plasma frequency, with the $2\omega_p$ peak approximately three orders of magnitude stronger than that in the s -polarized case [Fig. 4(a)]. This enhanced response is consistent with efficient excitation of bulk plasma waves driven by vacuum-heated electrons, the dominant collisionless absorption mechanism in this subrelativistic regime [68–70].

By increasing the laser strength to the relativistic regime ($a_0 = 1$), the electron plasma frequency is modified to $\omega_p = \sqrt{n_e e^2 / \gamma m_e \epsilon_0}$ and the corresponding $2\omega_L$ critical density becomes $n_{c(2\omega_L)} = 4\gamma n_c$, where $\gamma = (1 + a_0^2/2)^{1/2}$. Higher a_0 values thus lead to deeper laser penetration. At the relativistic critical surface, dense electron jets become more pronounced, exhibiting excursions up to $\sim 0.2\lambda_L$ in length [Fig. 3(h)], resembling vacuum-heating-like motion that drives CWE harmonics [20]. The enhanced $\vec{J} \times \vec{B}$ force efficiently accelerates electrons into denser regions, as shown by the plasma oscillations throughout the density gradient in Fig. 3(h). The resulting current bunches shift slightly deeper into the gradient and become spatially elongated, with a

substantial fraction extending up to bulk density [Fig. 3(k)]. A structure at $2\omega_p$ is now observed, in contrast to Fig. 3(j), indicating strong electron heating. The harmonic spectrum extends to the plasma frequency before fading into the background, where one or two higher-order harmonics faintly emerge and the $2\omega_p$ emission is also visible [Fig. 4(e)]. In this relativistic regime, $\vec{J} \times \vec{B}$ heating is no longer negligible even in the p -polarized case, as evidenced by electron density jet oscillations occurring twice per optical cycle [Fig. 3(b)]. Similar current bunches are observed [Fig. 3(e)], yielding a harmonic spectrum of comparable amplitude [Fig. 4(b)] to the s -polarized case.

In the ultrarelativistic regime ($a_0 = 10$), significant plasma steepening leads to the loss of the density gradient. Density jets now emerge from regions within the bulk plasma and extend to the relativistic critical surface [Figs. 3(c) and 3(i)]. Electrons reentering the target traverse an effectively homogeneous density, suppressing harmonic generation via linear mode conversion [7]. The shift of the current bunches becomes pronounced, extending into the bulk and forming a planar structure oscillating beyond the plasma frequency [Figs. 3(f) and 3(l)]. Correspondingly, harmonics of appreciable normalized intensity extend beyond the plasma frequency, consistent with the ROM mechanism [Figs. 4(c) and 4(f)]. An odd-even asymmetry is observed in the s -polarized case, with the even harmonics displaying lower amplitudes. This asymmetry has been observed in previous simulation work under lower a_0 values, where it was more pronounced at the tail end of the spectrum [23]. Because the asymmetry is not observed experimentally, we attribute this feature to be simulation specific and resolving its origin would require a dedicated numerical study beyond the scope of the present work. The density oscillations and current bunches for p -polarization become indistinguishable from those driven by s -polarization. In the spectra, the odd harmonics exhibit amplitudes comparable to those driven by s -polarized interactions up to the 15th order, beyond which the p -polarized-driven harmonics exhibit greater amplitudes.

Despite neglecting multidimensional effects, the simulations qualitatively reproduce the spectral trends observed experimentally for s -polarized interaction: the presence of only few low-order harmonics at moderate intensities, consistent with prior work, a plasma-frequency cutoff at relativistic intensities, and harmonics extending beyond the cutoff in the ultrarelativistic regime, consistent with our observations. At relativistic intensities, simulations indicate that the harmonic spectra exhibit negligible dependence on incident polarization for orders well above the plasma-frequency cutoff. Experimentally, the peak photon yield of the 17th-order harmonic (23.5 nm) at UHI [Fig. 2(b)] reaches $\sim 1 \times 10^6$ detected photons, comparable to yields reported for p -polarized interactions at 45° and $a_0 = 30$ under optimized scale length conditions [62]. The difference in incidence angle and polarization between the two experiments reflects each interaction operating near a favorable geometry for its dominant absorption mechanism. At 45° , resonance and vacuum heating absorption preferentially enhance harmonic generation efficiency for p -polarized interactions [40], while at our shallower angle of 22.5° , the $\vec{J} \times \vec{B}$ mechanism dominates [23], enabling a meaningful comparison of the

photon yields observed across both geometries. Additional optimization methods enabling higher efficiencies for relativistic plasma harmonics are actively being explored [61]. The simulated differences in high-order harmonic amplitudes ($n > 15$) between s - and p -polarized cases remain consistent with prior work [46], which reports comparable amplitudes up to the 9th order for both polarizations. The discrepancy in a_0 between simulations and experiment can be attributed to the reduced plasma density used in simulations. For high-density targets, as in the experiment, overcoming the strong plasma restoring force necessitates correspondingly stronger driving fields to efficiently excite plasma oscillations. Additionally, extreme steepening toward the bulk as shown in Figs. 3(c) and 3(i) is unlikely; however, partial steepening can directly lead the laser to interact with relativistically oscillating electron jets, producing ROM harmonics, and enabling both mechanisms to coexist, as reflected in the two power scalings.

In conclusion, the experimental conditions isolate $\vec{J} \times \vec{B}$ heating as the dominant absorption mechanism, enabling harmonic generation across relativistic and ultrarelativistic intensities and revealing a transition from a pronounced plasma-frequency cutoff to harmonics extending beyond it. PIC simulations capture the evolution of harmonic generation

mechanisms driven solely by $\vec{J} \times \vec{B}$ heating, highlighting their strong intensity dependence. Understanding the role of $\vec{J} \times \vec{B}$ heating in the generation of CWE harmonics becomes particularly important for high- Z materials, which require higher laser intensities to reach fully ionized states and overcome the strong plasma restoring force. Exploiting this dominant absorption mechanism in future ultrarelativistic laser-solid experiments could suppress competing absorption pathways, thereby mitigating preplasma thermal expansion and enabling greater control over complex interactions, such as two-color-driven HHG.

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Data availability. The data that support the findings of this article are not publicly available upon publication because it is not technically feasible and/or the cost of preparing, depositing, and hosting the data would be prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

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