Experimental observation of acceleration-induced thermality

Morgan H. Lynch,[1,*] Eliahu Cohen,[2,†] Yaron Hadad,[1,‡] and Ido Kaminer[1,§]

[1]*Department of Electrical Engineering, Technion: Israel Institute of Technology, Haifa 32000, Israel*
[2]*Faculty of Engineering and the Institute of Nanotechnology and Advanced Materials, Bar Ilan University, Ramat Gan 5290002, Israel*

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We examine the radiation emitted by high-energy positrons channeled into silicon crystal samples. The positrons are modeled as semiclassical vector currents coupled to an Unruh-DeWitt detector to incorporate any local change in the energy of the positron. In the subsequent accelerated QED analysis, we discover a Larmor formula and power spectrum that are both thermalized by the acceleration. Thus, these systems explicitly exhibit thermalization of the detector energy gap at the celebrated Fulling-Davies-Unruh (FDU) temperature. Our derived power spectrum, with a nonzero energy gap, is then shown to have an excellent statistical agreement with high-energy channeling experiments and also provides a method to directly measure the FDU temperature. We also investigate the Rindler horizon dynamics and confirm that the Bekenstein-Hawking area-entropy law is satisfied in these experiments. As such, we present the evidence for the first observation of acceleration-induced thermality in a nonanalog system.

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I. INTRODUCTION

The machinery employed by quantum field theory in curved spacetime [1–5] to analyze radiation emission has the advantage of stripping away the minute details of the emission process and focusing entirely on the local change in energy of a radiating system. Motivated by the original analyses of Unruh [6] and DeWitt [7] to simplify the transitions in energy of quantum fields in general relativistic backgrounds by modeling them as two-level systems, we return to the original use of the "Unruh-DeWitt detector" to analyze how a radiating particle changes its energy. The fundamental insight from this perspective is that for certain deformations in the structure of spacetime an Unruh-DeWitt detector will undergo a transition, either up or down in energy, and radiate [6–12]. Remarkably, the character of this radiating system will be thermalized at a temperature defined by some characteristic inverse length scale of the spacetime or motion through it, e.g., acceleration, surface gravity, Hubble constant, or some combination thereof [6,13–16]. Each of these length scales will define the location of an event horizon, which, via quantum fluctuations near the horizon, will emit radiation which is thermalized at that characteristic temperature. In Friedmann-Robertson-Walker cosmologies [17], signatures of this temperature may be encoded in the anisotropies in the cosmic microwave background [18]. Analog systems involving "fluid black holes" [19–22] are also capable of exploring not only the temperature but other properties such as the entropy via correlations between Hawking pairs both inside and outside the effective horizon and even superradiant scattering. The most sought after temperature [23,24], due to the fact that it appears to be the most readily accessible experimentally, is the acceleration temperature $T = \frac{a}{2\pi}$. The path toward the discovery of this temperature began with the analysis of the quantum mechanical structure of the vacuum in inertial and accelerated reference frames by Fulling [25], by the flux of radiation from the $1 + 1$-dimensional moving mirror by Davies [26], and finally by the near horizon examination of Hawking radiation from black holes by Unruh [6]. Understanding this Fulling-Davies-Unruh (FDU) temperature has been the subject of a steadily growing community, as detailed in Ref. [23], and the techniques developed to explore it have spread to other fields and fostered their growth as well, e.g., the use of Unruh-DeWitt detectors in relativistic quantum information [27]. Since the FDU temperature and the general characteristics of electromagnetic radiation are determined by the acceleration, it is through accelerated electromagnetic systems that we expect to see the first signs of accelerated thermality.

These pursuits have culminated in an inherently thermodynamic understanding of the nature of relativistic quantum field theory in classical general relativistic backgrounds [28]. One of the first clues for this thermodynamic interpretation comes from the notion of black hole entropy. Bekenstein conjectured that the information content or entropy, S, associated with a black hole would be

*morgan.lynch@technion.ac.il
†eliahu.cohen@biu.ac.il
‡yaronhadad@gmail.com
§kaminer@technion.ac.il

proportional to the surface area, A, of the event horizon [29]. Hawking later determined the proportionality constant to be $1/4$ and thus gave rise to the Bekenstein-Hawking area-entropy law [13], $S = \frac{A}{4}$. What is particularly interesting about this expression is that it specifically depends on the Planck area, ℓ_p^2. This area-entropy law is applicable to more than just black hole and is, in fact, a general property of systems with horizons. In particular, the Rindler horizon associated with acceleration will also obey this law [30]. Consequently, given a system with sufficient acceleration, thermality can be verified and explored by both the presence of a well-defined FDU temperature and the Rindler horizon dynamics. In fact, they both provide an independent confirmation of the presence of acceleration-induced thermality.

Here, we employ a spacetime formulation of *accelerated quantum electrodynamics* (AQED) [8,9,31], via the use of a uniformly accelerated Unruh-DeWitt detector [6,7], and apply it to high-energy channeling radiation [32]. The result of this analysis is a statistically significant indication that the FDU temperature has finally been observed, in a nonanalog experiment. We demonstrate that a nonzero Unruh-DeWitt detector energy gap, which encodes local changes in energy of the radiating electron or positron, will not only be thermalized at the FDU temperature but also provide a significantly better explanation of high-energy channeling radiation than conventional models. We present the AQED response function and compute the power radiated by, and photon power spectrum of, a uniformly accelerated charge current. To compare with the channeling radiation experimental data, we assume an energy gap comprising a Taylor series in the photon frequency and an acceleration profile based on radiative energy loss. A chi-squared analysis is shown to strongly favor the presence of a nonzero energy gap and therefore provides strong evidence for the observation of thermality at the FDU temperature. We then provide an independent verification of thermality by analyzing the Rindler horizon dynamics. We find the Bekenstein-Hawking area-entropy law is satisfied via its convergence to $4\ell_p^2$. Here and throughout, we use natural units $\hbar = c = k_B = G = 1$.

II. AQED RESPONSE FUNCTION AND THERMALIZED OBSERVABLES

The calculation presented here uses an AQED approach for computing the power radiated from first principles. To this end, we make use of the current interaction [33], $\hat{S}_I = \int d^4x \hat{j}_\mu(x)\hat{A}^\mu(x)$, where $\hat{A}^\mu(x)$ is the second quantized photon field operator and $\hat{j}_\mu(x)$ is the electron current operator. If we couple an Unruh-DeWitt detector to the current, it takes the form $\hat{j}_\mu(x) = u_\mu \hat{q}(\tau)\delta^3(x - x_{\mathrm{tr}}(\tau))$. Here, u_μ is the 4-velocity of the electron, and $x_{\mathrm{tr}}(\tau)$ is the trajectory of the electron parametrized by its proper time. The monopole moment operator [7] is Heisenberg evolved via $\hat{q}(\tau) = e^{i\hat{H}\tau}\hat{q}(0)e^{-i\hat{H}\tau}$, and the charge of the electron is defined by $q = \langle E_f|\hat{q}(0)|E_i\rangle$, with $|E_i\rangle$ and $|E_f\rangle$ being the initial and final electron energy, respectively. We must also note that our monopole moment operator creates states of definite momentum. As such for energy gaps comprising a continuum of frequencies, such as that we will analyze in this work, our current should be thought of as coupled to a continuum of detectors, one for each frequency. This also will restrict our working regime to first order in perturbation theory. Then, by analyzing the amplitude, $\mathcal{A} = i\langle \mathbf{k}| \otimes \langle E_f|\hat{S}_i|E_i\rangle \otimes |0\rangle$, for the electron current to undergo a transition and emit one photon of momentum $\mathbf{k}$, we can compute the emission rate via the relativistic analog of Fermi's golden rule. As such, the AQED response function, see Sec. B of the Supplemental Material [34] and Ref. [1], is given by

$$\Gamma = q^2 \int d\xi e^{-i\Delta E\xi} U_{\mu\nu}[x', x] G^{\mu\nu}[x', x]. \tag{1}$$

Here, we see the standard Fourier transform of the Wightman function, $G^{\nu\mu}[x', x]$, but with indices which, for photons, represent the sum of polarization 4-vectors, $\sum_{i,j} \epsilon_i^\mu \epsilon_j^{*\nu}$, which contract with the 4-velocity product, $U_{\mu\nu} = u_\mu u_\nu$, to couple the motion to the allowed emission directions. This is nothing more than the byproduct of the standard $v \cdot A$ coupling that one typically sees in electron-photon systems. Recalling the Wightman function is given by the vacuum-to-vacuum two-point function $G^{\nu\mu}[x', x] = \langle 0|\hat{A}^{\dagger\nu}(x')\hat{A}^\mu(x)|0\rangle$, we will have

$$G^{\mu\nu}[x', x] = \frac{1}{(2\pi)^3}\frac{1}{2}\int \frac{d^3k}{\omega}\sum_{i,j}\epsilon_i^\mu \epsilon_j^{*\nu} e^{i(\mathbf{k}\cdot\Delta\mathbf{x} - \omega(t'-t))}. \tag{2}$$

Since we wish to compute the power radiated, $\mathcal{S}$, we must also include an additional factor of frequency in the above Wightman function [9,11,12], as $\mathcal{S} = \int \frac{d\Gamma}{d\omega}\omega d\omega$. When we compute the power radiated by an accelerated electron, we should expect to recover the Larmor formula. As we shall see, this is indeed the case. For a trajectory with constant proper acceleration a, parametrized by the proper time τ, we will have hyperbolic motion with 4-velocity $u^\mu = (\cosh(a\tau), 0, 0, \sinh(a\tau))$ [23]. Using this trajectory, we arrive at a rather striking result. It is expected that we should obtain the Larmor formula, but there is an intermediate result: the fact we have uniformly accelerated motion implies that we should also find signatures of thermality. Explicit computation of the power, see Sec. C of the Supplemental Material, yields

$$\mathcal{S} = \frac{2}{3}\alpha a^2 \frac{1}{1 + e^{2\pi\Delta E/a}}. \tag{3}$$

It is indeed surprising that we have obtained the Larmor formula that is thermalized at the FDU temperature.

We must also point out the change in statistics from bosonic to fermionic. This is characteristic of accelerated thermal observables where the statistics depends on, for example, powers of frequency ω in the computation, number of particles emitted, and/or the dimensionality of the system [8,9,11,35]. To recover the standard Larmor formula, we must compute the total power by summing over emission and absorption of zero energy Rindler photons [24,36–38], i.e., $\lim_{\Delta E\to 0}\mathcal{S}(\Delta E)+\mathcal{S}(-\Delta E)=\frac{2}{3}\alpha a^2$. This is also corroborated by the notion that classical radiation sources can be viewed as "gapless" limit of an Unruh-DeWitt detector [39]. To compare the AQED theory with the experimental data, we also present the total power radiated per unit frequency; see the Supplemental Material, Eqs. (S41) and (S42),

$$\frac{d\mathcal{S}}{d\omega}=-i\frac{4}{3}\alpha\frac{\omega^2}{a}\left[\delta H^{(2)}_{\frac{2i\Delta E}{a}}\left(-\frac{2i\omega\gamma}{a}\right)-\frac{1}{2}\left(H^{(2)}_{\frac{2i\Delta E}{a}-2}\left(-\frac{2i\omega\gamma}{a}\right)+H^{(2)}_{\frac{2i\Delta E}{a}+2}\left(-\frac{2i\omega\gamma}{a}\right)\right)\right][1+e^{2\pi\Delta E/a}]. \tag{4}$$

Here, we have defined the relativistic boost parameter $\delta=2\gamma^2-1$ and see that the power spectrum comprises the Hankel function of the second kind, $H^{(2)}_{\ell}(x)$ [40]. We also made the presence of thermality more apparent by making use of the following identity: $H^{(2)}_{\ell}(x)=e^{i\ell\pi}H^{(2)}_{-\ell}(x)$. The implication of this property of Hankel functions of the second kind is the manifest detailed balance at thermal equilibrium of the power spectrum by rigorous mathematical identity. The exponent produced by the change in sign of the Hankel index is precisely the Boltzmann factor comprising the Unruh-DeWitt detector energy gap thermalized at the temperature $T=\frac{a}{2\pi}$. As such, we are led to the conclusion that systems described by this power spectrum imply the experimental observations of thermality at the FDU temperature.

This thermal phenomenon, commonly referred to as the Unruh effect [6], has had a considerable amount of effort dedicated to its study as well as a detailed exploration of potential experimental settings which could measure it [23]. The main difficulty with measuring such an effect is the vanishingly small energy scale set by the acceleration when compared to the scale of the energy gap. Broadly speaking, the ability to probe the Unruh effect necessitates $|\Delta E|\sim\frac{a}{2\pi}$ so that the thermal distribution can be explored. The difficulty then lies in bringing the two energy scales together: finding acceleration scales that reach the energy gap from below or energy gaps that can reach the small acceleration scale from above. If both a small energy gap and a large acceleration scale can be found in an experimental system, then this would provide the best chance of measuring the Unruh effect. Through this logic, high acceleration scales coupled to small energy gaps, we apply the above power spectrum, Eq. (4), to the recent channeling radiation experiment in aligned crystals [32]. There, we have LHC scale energetic positrons channeled into single crystal silicon (large acceleration) sensitive to the channeling oscillation [41], recoil/radiation reaction [42–46], and other processes in the comoving frame (small energy gap). Note that the use of Unruh-DeWitt detectors has also been successfully applied to exploring recoil/radiation reaction in accelerated Cherenkov systems [31]. As we shall see, it appears to be the case that channeling radiation not only provides a setting to measure radiation reaction but may finally enable a system to explore the Unruh effect experimentally.

III. INCORPORATING RADIATION REACTION

Understanding the nature of recoil due to photon emission from an accelerated charge has been one of the longest-standing problems in physics: the problem of radiation reaction. The problem itself is typically examined by including recoil, based on the Larmor formula, in the Lorentz force, yielding the Lorentz-Abraham-Dirac (LAD) equation [42]. However, this formalism is plagued by runaway solutions which have yet to be tamed. Approximations to this equation were soon discovered [43] and subsequently solved [44,45]. Despite this progress, the problem of runaway solutions still persists [46]. Perhaps, it may be easier to incorporate recoil not at the level of the Lorentz force but rather incorporate it into the computation of an observable using an Unruh-DeWitt detector.

To get a better understanding of how to incorporate recoil, let us first examine a much simpler scenario, that of Cherenkov radiation. The reason for this is because the quantum recoil, or radiation reaction, correction to Cherenkov radiation is already well known. When we include an Unruh-DeWitt detector into the Cherenkov regime, we obtain the following Frank-Tamm formula [31]:

$$\frac{d\Gamma}{d\omega}=\alpha\beta\left[1-\left[\frac{1}{n\beta}+\frac{\Delta E}{n\beta\omega\gamma}\right]^2\right]. \tag{5}$$

The presence of the energy gap in the expression yields the anomalous Doppler effect [47]. The interesting thing is that when we set the energy gap to $\Delta E=\frac{(n^2-1)\omega^2}{2m}$ then we reproduce the quantum recoil correction to Cherenkov radiation identically [48]. The next question is where our energy gap, which gave us the recoil correction, came from. If we consider an electron with a renormalized rest mass [49] comprising the electron rest mass and the emitted photons energy $E_i=\sqrt{m^2+\omega^2}$ and a final-state electron energy comprising a bare electron and recoil momentum $k=n\omega$ produced by the photon emission, $E_f=\sqrt{m^2+(n\omega)^2}$, then the difference in energy is given by $\Delta E=E_f-E_i\approx\frac{(n^2-1)\omega^2}{2m}$. It is this physical interpretation, backed by mass renormalization in the Cherenkov regime, that we will utilize in the Larmor setting for the inclusion of recoil, i.e., radiation reaction.

Having successfully formulated a theory of recoil in Cherenkov emission using an Unruh-DeWitt detector, we now turn to the inclusion of recoil into the Larmor formula. Given that we know the inclusion of recoil is tantamount to including a recoil kinetic energy, which is process independent, we then look for a complete derivation of the Larmor formula, within an Unruh framework, and include an energy gap $\Delta E \sim \frac{\omega^2}{2m}$. This form of recoil is also present when looking at nonrelativistic Unruh-DeWitt detectors with a quantized center of mass [50]. The incorporation of a quantized center of mass also gives rise to the possibility of exploring light matter interactions which go beyond the Unruh-DeWitt detector approximation [51]. By incorporating this recoil into our response function, we can compute, via a series expansion in ω, the total power radiated. The first term gives the Larmor formula, $\mathcal{S}_0 = \frac{2}{3}\alpha a^2$. The second term gives the 1st order quantum recoil correction, $\mathcal{S}_1 = -\frac{8\alpha}{m} a^2 T_{\mathrm{FDU}}$; see Sec. D of the Supplemental Material. Combining the two gives us our quantum corrected Larmor formula,

$$\mathcal{S} = \frac{2}{3}\alpha a^2 \left[1 - \frac{12}{m} T_{\mathrm{FDU}}\right]. \tag{6}$$

The functional dependence on the temperature as well as the sign are also in complete agreement with the first order correction to the massless scalar case [52]. We can also compute the radiation reaction force for each term. If we recall that the work done by the radiation reaction force is equal to the total energy radiated, we have $\int F^{\mathrm{rr}} dx = -\int \mathcal{S} dt$. From this, we find our radiation reaction forces to be $F_0^{\mathrm{rr}} = \frac{2}{3}\alpha J$ and $F_1^{\mathrm{rr}} = -\frac{8\alpha}{\pi m} J a$. Thus, we find the first-order quantum LAD equation to be

$$m\frac{du^\mu}{ds} = qF^{\mu\nu}u_\nu + \frac{2}{3}\alpha\left[1 - \frac{24}{m}T_{\mathrm{FDU}}\right][J^\mu + a^2 u^\mu]. \tag{7}$$

It is interesting to note that the quantum recoil correction seems to renormalize the fine structure constant by the acceleration $\alpha \to \alpha[1 - \frac{24}{m} T_{\mathrm{FDU}}]$. It is with the above considerations, namely, the inclusion of $\Delta E \sim \frac{\omega^2}{2m}$ in our energy gap, that we may investigate recoil or radiation reaction in accelerated systems; see Sec. J of the Supplemental Material for more detail. The existence of such a term in the energy gap of our power spectrum, Eq. (4), when compared to the channeling radiation data set [32], would indeed confirm the presence of recoil.

IV. MEASURING THE FULLING-DAVIES-UNRUH TEMPERATURE

A recent proposal was made by Cozzella *et al.*, which outlined a method of measuring the FDU temperature directly from a dataset [24]. In their analysis, a longitudinally accelerated electron is subjected to a comoving cyclotron oscillation in the transverse plane, and they compute the photon emission spectrum per unit transverse momentum. The analysis was carried out in both Minkowski space, i.e., the laboratory frame, and Rindler space, i.e., the comoving frame. General covariance necessitates that the emission rates be identical, and in order to accomplish this, the Rindler frame computation must take into account the emission and absorption of Rindler photons from a background thermal distribution at the FDU temperature; In other words, the Unruh effect is mandatory to render the process covariant [12]. Most importantly, the temperature of the Rindler bath can be left arbitrary, and the resultant observable can be matched to the dataset by finding the best-fit temperature, thereby providing a method to directly measure the FDU temperature.

What is particularly intriguing about this proposal is that it closely resembles the dynamics of the channeling radiation experiment [32]. There, a positron undergoes a longitudinal acceleration and oscillates transversely, i.e., a one-dimensional oscillation. The only real difference is the fact that the channeling experiment is ultrarelativistic. As such, if we are able to reproduce our thermalized power spectrum via a Rindler frame computation, we cannot only gain insight into the nature of the processes present in the comoving frame but can then also leave the Rindler temperature term arbitrary to yield a more general expression to be applied to the channeling dataset. Should we find that the power spectrum accurately describes the data, we could then provide a direct measurement of the temperature, thereby confirming the presence of a thermalized Rindler bath and thus the Unruh effect itself.

To analyze the Rindler space emission rate, we will utilize the formalism developed in Refs. [24,53], and we must first consider the Rindler coordinate system (τ, ξ, x, y), with the line element given by $ds^2 = e^{2a\xi}(d\tau^2 - d\xi^2) - dx_\perp^2$. The two Rindler coordinates (τ, ξ) are related to the laboratory time and z coordinate via $t = (e^{a\xi}/a)\sinh(a\tau)$ and $z = (e^{a\xi}/a)\cosh(a\tau)$. To describe the channeling oscillation, we adopt a nonrelativistic comoving transverse oscillation, in the Rindler coordinate system, described by the 4-velocity $u^\mu = (1, 0, v_0 \cos(\Omega\tau), 0)$. Here, the amplitude of the channeling oscillation velocity $v_0 = A\Omega$, with A being the amplitude of the oscillation and Ω the channeling oscillation frequency in the comoving frame.

It has been well established that photon emission into Minkowksi space corresponds to both the emission of Rindler photons from and absorption of Rindler photons into the thermalized Rindler bath, due to the Unruh effect, in the comoving frame [23,24,36,53]. This is accomplished by weighting the probability of absorption of a Rindler photon by a thermal distribution and the emission probability by a thermal distribution plus 1, i.e., $\mathcal{P}_{abs} \sim |\mathcal{A}_{abs}|^2(1/(e^{\omega_r/T}-1))$ and $\mathcal{P}_{\rm emi} \sim |\mathcal{A}_{abs}|^2(1+1/(e^{\omega_r/T}-1))$. Note this also makes use of the fact that $|\mathcal{A}_{abs}|^2 = |\mathcal{A}_{\rm emi}|^2$. The temperature of the background thermal bath, T, is kept arbitrary. The total Rindler emission rate is then given by, see Sec. G of the Supplemental Material,

$$\Gamma_r = \int_{\infty}^{\infty} d^2k_\perp \int_0^\infty d\omega_r \frac{[|\mathcal{A}^1_{abs}|^2 + |\mathcal{A}^2_{abs}|^2]}{\Delta\tau} \coth(\omega_r/(2T)). \tag{8}$$

Here, the two terms correspond to summing over both Rindler photon polarizations. Each polarization is described by the mode functions, $f(x) = K_{i\omega_r/a}(\frac{k_\perp}{a}e^{a\xi}) \times e^{i(k_\perp \cdot x_\perp - \omega_r\tau)}$. Note that these photons comprise plane waves transverse to the direction of acceleration and a K2 modified Bessel function of the second kind function along the acceleration axis. The transverse plane waves are characterized by their momenta $k_\perp$, and the K2 Bessel is characterized by the Rindler frequency ω_r. What is important to note at this point is that the contribution of the background thermal distribution is contained in the factor, $\coth(\omega_r/(2T))$. When analyzing the case of Rindler photon emission and absorption due to our transverse channeling oscillation, we then obtain the following total emission spectrum per transverse momentum:

$$\frac{d\Gamma_r}{d^2k_\perp} = -\frac{i\alpha}{2\pi a}\sinh(\pi\omega_r/a)\coth(\omega_r/(2T))e^{\frac{\pi\omega_r}{a}} \times \left[\delta H^{(2)}_{\frac{2i\omega_r}{a}}\left(\frac{-2ik_\perp\gamma}{a}\right) - \frac{1}{2}\left(H^{(2)}_{\frac{2i\omega_r}{a}+2}\left(\frac{-2ik_\perp\gamma}{a}\right) + H^{(2)}_{\frac{2i\omega_r}{a}-2}\left(\frac{-2ik_\perp\gamma}{a}\right)\right)\right]. \tag{9}$$

Here, we see that both the thermal contribution as well as the indices of the Hankel functions are determined by ω_r. For our transverse channeling oscillation, this sets $\omega_r = k_x A\Omega$. We must also comment on the fact that we have two new thermal terms. First, we have an additional factor of $\sinh(\pi\omega_r/a)$ due to the Rindler mode, $K_n(x)$, normalization. Second, we have a factor $e^{\frac{\pi\omega_r}{a}}$ which comes from the mode transformation $K_{i\omega_r/a}(k_\perp/a) \sim e^{\omega_r/T_{\rm FDU}} H_{i\omega_r/a}(-ik_\perp/a)$. Now, by comparing this expression to that of the case of Minkowski photon emission, we verify that the expressions match when the temperature of the Rindler bath is set to the FDU temperature. When analyzing photon emission using an Unruh-DeWitt detector in Minkowski space, we must sum over both the excitation and deexcitation of the detector [39], i.e., $\frac{d\Gamma_{\rm tot}}{d^2k_\perp} = \frac{d\Gamma_{\Delta E}}{d^2k_\perp} + \frac{d\Gamma_{-\Delta E}}{d^2k_\perp}$. The resulting Minkowski emission rate yields

$$\Gamma_m = \frac{1}{\Delta\tau}\int d^2k_\perp dk_z[|A^m_\uparrow|^2 + |A^m_\downarrow|^2] = \frac{1}{\Delta\tau}\int d^2k_\perp dk_z|A^m_\uparrow|^2[1+e^{\Delta E/T_{\rm FDU}}]. \tag{10}$$

Note that here we made use of the detailed balance relationship, $\Gamma^m_\downarrow = \Gamma^m_\uparrow e^{\Delta E/T_{\rm FDU}}$, demonstrating thermal equilibrium at the FDU temperature. What is important to note is that the presence of thermality is entirely encoded in the Boltzmann-like factor $1+e^{\Delta E/T_{\rm FDU}}$ and does not depend on the process specific Minkowski amplitudes A^m. Of course, we know this relation is also manifested in the emission rate via our Hankel modes, $H^{(2)}_{\frac{-2i\Delta E}{a}}(x) = e^{2\pi\Delta E/a}H^{(2)}_{\frac{2i\Delta E}{a}}(x)$. As such, our transverse Minkowski emission spectrum is given by

$$\frac{d\Gamma_m}{d^2k_\perp} = \frac{-i\alpha}{4\pi a}\left[\delta H^{(2)}_{\frac{2i\Delta E}{a}}\left(-\frac{2ik_\perp\gamma}{a}\right) - \frac{1}{2}\left(H^{(2)}_{\frac{2i\Delta E}{a}-2}\left(-\frac{2ik_\perp\gamma}{a}\right) + H^{(2)}_{\frac{2i\Delta E}{a}+2}\left(-\frac{2ik_\perp\gamma}{a}\right)\right)\right][1+e^{2\pi\Delta E/a}]. \tag{11}$$

Now, given the fact that both Minkowski and Rindler emission rates describe precisely the same physics, the rates are then equal $\frac{d\Gamma_m}{d^2k_\perp} = \frac{d\Gamma_r}{d^2k_\perp}$. This then yields our thermal transformation function,

$$[1+e^{\Delta E/T_{\rm FDU}}] = \frac{1}{2}\sinh(\pi\omega_r/a)\coth(\omega_r/(2T))e^{\frac{\pi\omega_r}{a}}. \tag{12}$$

We see that the above emission spectra, computed in both Rindler space with a background thermal bath and in vacuum Minkowski space, are indeed identical, provided we identify our energy gap with Rindler frequency $\Delta E = \omega_r$ and the background temperature with the FDU temperature $T = \frac{a}{2\pi}$. Functionally, what we have done is utilize the covariance of the photon emission between the Minkowski and Rindler frames in order to "track" the presence of thermality. We have detailed balance from the Minkowski side, along with the photon exchange with the Rindler bath on the Rindler side scaled by the mode normalization. This relationship is process independent and demonstrates the fundamental nature of how the Unruh effect is mandatory for the self-consistency of quantum field theory [12]; detailed balance in Minkowski space corresponds to photon exchange with Rindler bath in Rindler space.

Now, based on the thermal identity, Eq. (12), we can write our full Minkowski power spectrum with the temperature of the Rindler bath left arbitrary. In short, we are converting the detailed balance relation back into the Rindler thermal bath. This will allow us to directly measure the temperature from the dataset. Hence,

$$\frac{dS}{d\omega} = -i\frac{4}{3}\alpha\frac{\omega^2}{a}\left[\delta H^{(2)}_{\frac{2i\Delta E}{a}}\left(-\frac{2i\omega\gamma}{a}\right) - \frac{1}{2}\left(H^{(2)}_{\frac{2i\Delta E}{a}-2}\left(-\frac{2i\omega\gamma}{a}\right) + H^{(2)}_{\frac{2i\Delta E}{a}+2}\left(-\frac{2i\omega\gamma}{a}\right)\right)\right] \times \sinh(\pi\Delta E/a)\coth(\Delta E/(2T))e^{\frac{\pi\Delta E}{a}}. \quad (13)$$

Provided the thermalized power spectrum, Eq. (4), accurately describes the channeling data, we can then use the above expression to fit the temperature. A successful measurement of the FDU temperature would confirm the presence of a thermal bath in the Rindler frame and thus the Unruh effect, thereby realizing the proposal of Cozzella *et al.* [24].

V. BEKENSTEIN-HAWKING AREA-ENTROPY LAW

Given a system which is manifestly thermalized by the acceleration also prompts an additional analysis of the horizon thermodynamics. Building upon the work of Bekenstein, who reasoned that the entropy of a black hole would be proportional to the area, Hawking, upon the discovery of black hole evaporation, was able to fix the proportionality constant to $1/4$. This gave birth to the Bekenstein-Hawking entropy-area law [13,29], $S = \frac{A}{4}$. Measuring such a quantity from an astrophysical black hole seems well beyond our current experimental capabilities. However, systems which are thermalized by the acceleration obey the same area-entropy law due to the Unruh effect [30]. As such, for highly accelerated systems, one can directly test the hypothesis of Bekenstein and Hawking.

In the proper frame, the change in the horizon area is determined by the amount of energy radiated by each positron, $\Delta\tilde{E}$, into the horizon. This quantity we will extract from the actual data via $\Delta\tilde{E} = \frac{4c}{3x_0}\int\frac{dS_{\mathrm{data}}}{d\omega}d\omega dt$; see Sec. A of the Supplemental Material for more details. The corresponding change in horizon area [30] is given by $\Delta A = \frac{Gc^5}{\hbar}\frac{8\pi m^3\Delta\tilde{E}}{E_i^3 a}$; see Sec. I of the Supplemental Material. From the first law of thermodynamics, we can also determine the entropy difference in the proper frame based on the difference in positron energy in the laboratory frame, $\Delta S = \frac{c^8}{\hbar^2}\frac{\pi m^3}{a}\left[\frac{1}{(E_i - \Delta\tilde{E})^2} - \frac{1}{E_i^2}\right]$. Note that we reintroduced all fundamental constants into the above expressions for the area and entropy change so we can resolve the presence of the Planck area, $\ell_p^2 = \frac{G\hbar}{c^3}$. Finally, if the hypothesis of Bekenstein and Hawking is true, then we will have $\frac{\Delta A}{\Delta S} = 4\ell_p^2$. As such, we will integrate the data in order to compute the following area to entropy ratio, see Sec. I of the Supplemental Material,

$$\frac{\Delta A}{\Delta S} = \ell_p^2\frac{8\Delta\tilde{E}}{E_i^3}\left[\frac{1}{(E_i - \Delta\tilde{E})^2} - \frac{1}{E_i^2}\right]^{-1}. \quad (14)$$

When the system thermalizes, it is expected that the above expression, with the appropriate integration over the data, will converge to $4\ell_p^2$. What is required to confirm the presence of thermality is that this expression satisfies the relation $S = A/4$, even with $\Delta E \neq 0$. Conceptually, what this means is that by the original integration of the first law, we fix $S_i = \frac{A_i}{4}$. This is the zeroth-order "initial condition" of the integration. Then, we must have ΔE evolve in such a way that the change in the area and entropy also obey $\Delta S = \frac{\Delta A}{4}$. This is not always the case in fact. The ΔE that presents itself in the above expressions must come from a thermalized observable. When applied to the high-energy channeling experiment, a verification of this area-entropy law will provide independent corroborating evidence for the presence of acceleration-induced thermality which would compliment the analysis based on the Minkowski and Rindler thermal power spectra.

VI. EXPERIMENTAL OBSERVATION OF ACCELERATION-INDUCED THERMALITY IN CHANNELING RADIATION

When a highly energetic charged particle propagates in a material, it will lose energy and emit a smooth spectrum of radiation known as bremsstrahlung [54]. This particular process is associated with scattering off individual atomic sites that have a random distribution, i.e., an amorphous crystalline structure for solids or the random distribution of atoms in the case of a liquid or gas. However, if we have a solid with some periodic crystalline structure, then we can look at motion along an axis of symmetry where the charged particle is "channeled" and will propagate in an effective hollow wave guide, or two-dimensional (2D) plane between atomic layers, produced by the structure. Then, the charged particle is confined to the potential well and can oscillate back and forth transversely to its direction of propagation and therefore radiate. This process is known as channeling radiation [55,56]. In the channeling experiment [32], the data from the photon power spectrum produced by the rapidly decelerated positrons are presented and are precisely the observable computed in our analysis. There, 178.2 GeV positrons are fired into a 3.8 mm sample of single crystal silicon aligned along the $\langle 111\rangle$ axis. Because the positron is positively charged, it will be repelled by the positively charged atomic sites and therefore be subject to a confining harmonic oscillator potential. The transitions between these harmonic oscillator states give rise to the channeling radiation seen in the experiment. These transitions between states, i.e., the transverse oscillation, decouple from the longitudinal motion, i.e., the beam velocity. The energy loss produced by this photon emission then gives rise to the large accelerations necessary to bring about thermality.

If we consider a characteristic laboratory frame photon frequency, ω_0, produced by the channeling radiation, we can determine the momentum scale of the associated process. This will result in a change in momentum imparted on the positron given by $|\Delta p| = |k_0| = \omega_0$. This momentum change occurs during the timescale of the emission process. Taking the physical size of the photon to be $\Delta x = \lambda_0/2$, we can then determine the emission time to be $\Delta t = \Delta x/c = \frac{\pi}{\omega_0}$. Then, recall that the relativistic version of Newton's law, $f = \frac{\Delta p}{\Delta t} = ma'$, will enable us to determine the proper acceleration, a'; see Sec. H of the Supplemental Material. Hence, we find our proper acceleration is given by

$$a' = \frac{\omega_0^2}{\pi m} \tag{15}$$

This is a proper acceleration, but it is written in terms of the laboratory frequency. What is important to note is that when written in terms of proper quantities the proper acceleration, $a' = \frac{\omega'^2_{\max}\gamma^2}{\pi m}$, boosts as γ^2. More importantly, when computing the FDU temperature for the emission, we find a recoil/radiation reaction temperature, T_{RR}. Hence,

$$T_{\mathrm{RR}} = \frac{\omega_0'^2\gamma^2}{2m\pi^2}. \tag{16}$$

Note, we now have an FDU temperature which depends explicitly on the recoil kinetic energy, $\omega^2/2m$, which is imparted on the positron by the emission. It is this radiation reaction temperature, produced by the recoil itself, that we will look for in our experimental analysis.

To compare with the experimental data we need to convert all parameters to GeV [32,57]: $E_0 = 178.2$ GeV and $m = 0.000511$ GeV. We have an overall scale factor, s, to take into account detector efficiencies. To model the change in the positron's energy, we adopt a polynomial energy gap of the form $\Delta E = a_0 + a_1\omega + a_2\omega^2 + a_3\omega^3$; the logic being that the exact expression of the energy difference, produced by the relevant dispersion and conservation of momentum and energy relations associated with all processes present, will be amenable to a Taylor expansion in powers of the emitted photon's frequency [58]; see Sec. J of the Supplemental Material for more details. These parameters will be used to compute the total power spectrum, $\frac{dS_{\mathrm{tot}}}{d\omega} = \frac{dS(\Delta E)}{d\omega} + \frac{dS(-\Delta E)}{d\omega}$, with and without the energy gap, and then compared with the crystal data [32]. In principle, the a_0 term is associated with the channeling oscillation, a_1 is associated with the Rindler frequency, and the $a_2\omega^2$ term is associated with recoil or radiation reaction [42–46]. We included a frequency dependence up to an ω^3 term to capture any further dependence beyond the recoil term since that was what the experiment [32] measured. The energy gap can be "turned off" by setting it to zero. To better match the calculated spectrum to the data, we also include an overall scaling factor s to take into account detector efficiencies and other systematics of the experiment; see Sec. A of the Supplemental Material for more details. To ensure the presence of thermality, we must also examine the emission lifetime. Recalling the emission rate per unit frequency is given by $\frac{d\Gamma(\omega)}{d\omega} = \frac{dS(\omega)}{d\omega}\frac{1}{\omega}$, we then invert this and integrate up to a specific frequency to determine the thermalization time for that frequency, i.e., the partial thermalization time [59], $t(\omega) = \frac{1}{\int_0^\omega \frac{d\Gamma(\omega')}{d\omega'} d\omega'}$. Thermality necessitates that we must have this lifetime be shorter than the travel time within the crystal. This will then require a low-frequency cutoff for frequencies that do not have enough time for thermalization to take hold. We then the compute best fits of our theory, with respect to the paramaters s, Fulling-Davies-Unruh temperature T_{FDU}, Rindler temperature T_R, and a_i, along with the reduced chi-squared statistic for each low-frequency cutoff.

The plot of our emission lifetime and reduced chi-squared statistics for each cutoff for the 3.8 mm channeling crystal [32] is presented in Fig. 1. Also included is the best-fit power spectra (4) along with examples of changing the Rindler bath temperature (13) and data [32], which yielded the first statistically significant signals of accelerated thermality based on the chi-squared statistic. To verify the presence of thermality, we also examine the Bekenstein-Hawking area-entropy law (14). It is through its convergence to $4\ell_p^2$, along with the chi squared of order unity, that we are led to the conclusion of acceleration-induced thermality.

To summarize our analysis, we found sufficient time inside the crystal for the system to be thermalized by the acceleration. The thermalization time implied a low-energy cutoff of approximately 22 GeV for the channeling data. Performing best fits of our theory to the data, with cutoffs at multiples of 10 GeV, yielded a reduced chi-squared statistic within the 1 standard deviation threshold at 30 GeV. For this cutoff, the overall chi-squared statistic also favors a nonzero energy gap with differences, $\Delta\chi^2 = |\chi^2_{\mathrm{gap}} - \chi^2_{\mathrm{no\,gap}}| = 364$. The best-fit parameter for the ω^2 term is $a_2 = 846\ \mathrm{GeV}^{-1}$, which accurately reflects the expected value of $\frac{1}{2m} \sim 978.2\ \mathrm{GeV}^{-1}$, thereby confirming the presence of recoil and thus radiation reaction. The best-fit value for channeling frequency is $a_0 = 5.6$ eV, in the proper frame, which is typical for channeling oscillations [56]. We can also confirm the presence of the fiducial Rindler frequency term, $\omega_r \sim \omega A\Omega$, by the coefficient on the linear term ω. For an amplitude on the order of the lattice constant for silicon, $A = 5.4$ Å, and using the best fit value for the channeling frequency, $\Omega = a_0$, we expect the linear coefficient to be $A\Omega \sim .015$. We indeed find our best-fit value to be $a_1 = .012$, thereby confirming the Rindler frame analysis, i.e., that Minkowski photon emission corresponds, at the very least, to the emission and absorption of Rindler photons at the Rindler frequency set

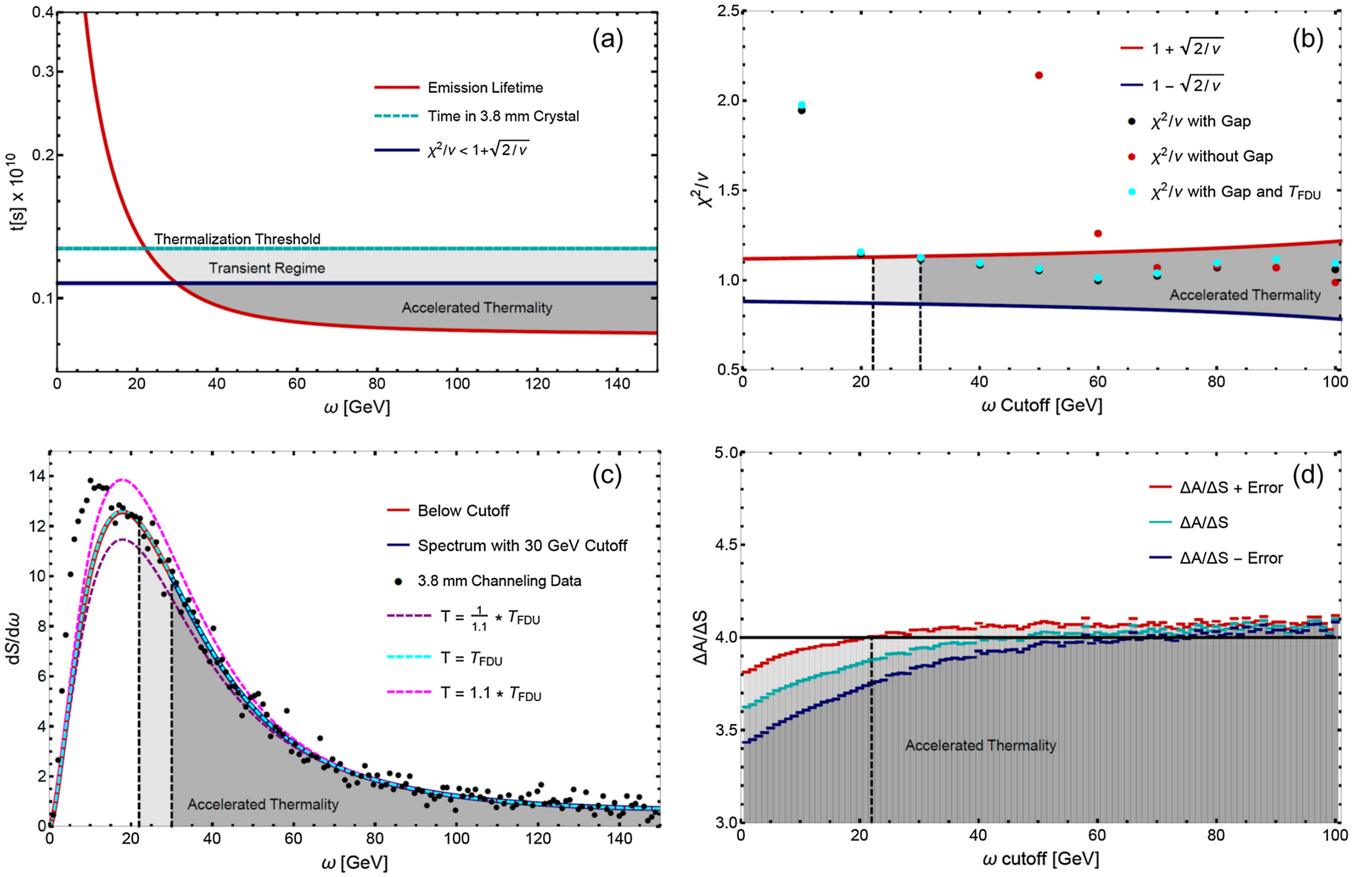

FIG. 1. Experimental observation of acceleration-induced thermality. For the Unruh effect to manifest itself, whatever process that takes place must have time to thermalize. Plot (a) demonstrates that the emission lifetime of the channeling experiment is shorter than the time the positron traverses the crystal sample, thereby demonstrating that the system has time to thermalize. Note the threshold energy is approximately 22 GeV. With the threshold of thermality exceeded, we must then look at various cutoffs beyond this threshold for a statistical signal which will confirm it. Plot (b) shows the chi-squared statistic for the best-fit power spectrum with various low-frequency cutoffs. Included is also the chi squared for the power spectrum, Eq. (13), with the best-fit Rindler temperatures. The channeling radiation has a $\chi^2/\nu < 1 + \sqrt{2/\nu}$ starting at 30 GeV. Moreover, the presence of a thermalized energy gap is favored over no gap. Plot (c) presents the power spectrum with a thermalized energy gap along with the channeling data for the first cutoff, 30 GeV, that satisfies the chi-squared criterion. Also included is the same power spectrum with the Rindler temperature set to T_{FDU}, $(1.1)T_{\mathrm{FDU}}$, and $(1/1.1)T_{\mathrm{FDU}}$ to illustrate the sensitivity of the spectrum to the temperature. The thermalized power spectrum (4), (13) shows excellent agreement with the dataset and provides evidence for the first observation of the Unruh effect. Plot (d) shows the ratio of the Rindler horizon area to the entropy as a function of the low-energy cutoff. For thermalized systems, this ratio is identically $4\ell_p^2$. Again, note the thermalization threshold energy of approximately 22 GeV, thereby providing corroborating evidence for acceleration-induced thermality.

by the channeling oscillation $\omega_r \sim A\Omega$, [24,53]. This implies that our energy gap is given by $\Delta E \sim \Omega + \Omega A\omega + \frac{\omega^2}{2m}$: the sum of a pure channeling frequency Ω, the Rindler frequency $\omega_r = A\Omega\omega$, and the recoil term $\frac{\omega^2}{2m}$.

By keeping the FDU temperature arbitrary and performing best fits for the acceleration for each power spectrum, we can measure the FDU temperature, T_{FDU}, with the eight cutoffs, which satisfy the chi-squared criterion: 30–100 GeV. Then, by fixing the acceleration to the best-fit value for each cutoff, we can perform the best fit for the Rindler temperature, T_R, specifically for these same cutoffs. Note the source of error for both these measurements is the standard deviation from the mean and therefore will decrease as $1/\sqrt{n}$, which for our current analysis utilized $n = 8$. For the characteristic frequency of the radiation reaction temperature, T_{RR}, let us use the first frequency to thermalize from the dataset, i.e., $\omega_0 = 150$ GeV. Note that this will also serve as an upper bound on the temperature. We can now compare the experimental measurements for the average FDU temperature, the average Rindler temperature, and the radiation reaction temperature. Hence,

$$\begin{aligned} T_{\mathrm{FDU}} &= 1.80 \pm .51 \text{ PeV} \\ T_R &= 1.96 \pm .49 \text{ PeV} \\ T_{\mathrm{RR}} &= 2.23 \text{ PeV}. \end{aligned} \tag{17}$$

As such, we find the temperature of the Rindler bath to be $T_R = T_{\mathrm{FDU}}(1.09 \pm .41)$. This experimental measurement confirms the predictions of Fulling [25], Davies [26], and Unruh [6] and also realizes the proposal put forward by Cozzella *et al.* for confirming the presence of the Unruh effect by directly measuring the FDU temperature [24]. The range of frequencies which set the radiation reaction temperature is also given by $\omega_0 = 135 \pm 17.6$ GeV. The agreement of both the FDU temperature with the radiation reaction temperature also corroborates the presence of recoil. This result seems natural under framework of radiation reaction; the acceleration scale of the system appears to be set by the recoil energy of the first frequencies to thermalize, in this case $\omega_0 \sim 150$ GeV. This immense recoil acceleration, $a \sim 5 \times 10^{39}\ m/s^2$, gives rise to a temperature of 2×10^{19} K. The incredibly high-energy scales here are due to the fact that for the channeling experiment [32] the Lorentz gamma is $\gamma = 3.5 \times 10^5$. Other prospects for future experiments could even include two-photon emission processes [60] as well as explore finite-size effects such as gauge invariance in minimal and dipole coupling interactions, [51,61].

In short, the first three terms in the Unruh-DeWitt detector energy gap all correspond to expected values. In particular, the presence of an approximately $\frac{\omega^2}{2m}$ term confirms the presence of recoil. The theory employed in the analysis obeys detailed balance, is thermalized at the FDU temperature, and also reproduces the Larmor formula in the appropriate limit [24,36–38]. Through a complimentary Rindler analysis, we confirmed the presence of a thermal bath of Rindler photons at the FDU temperature in the comoving frame [24,53]. Then, through a completely independent thermodynamic analysis, we found that the Bekenstein-Hawking area-entropy law was also satisfied starting at the same approximately 22 GeV threshold. It is through the chi-squared analysis of the best-fit power spectrum, measurement of the FDU temperature, and confirmation of the Bekenstein-Hawking area-entropy law that we are led to the conclusion of acceleration-induced thermality.

VII. CONCLUSIONS

The main focus of this manuscript deals with the presence of accelerated thermality in a high-energy channeling radiation experiment that measured radiation reaction [32]. However, there are two additional experiments that also report evidence of radiation reaction using laser-wakefield acceleration [62,63]. There, rather than the photon power spectrum, the final state electron energy was measured. It would be an interesting avenue of research to apply this formalism to their experiments in search of accelerated thermality as well. The connection between the Unruh effect and radiation reaction has long been discussed in the literature [64,65], and it appears that these, and future, systems may provide a robust experimental setting to not only explore these intriguing aspects of radiation reaction but also to investigate the intriguing nature of acceleration-induced thermality.

In this article, we employed the theory of accelerated quantum electrodynamics and used it to explore the radiation produced by uniform accelerated motion. When applied to the problem of channeling radiation, we are able to incorporate a local change in energy, utilizing techniques from quantum field theory in curved spacetime, by setting the energy gap of an Unruh-DeWitt detector equal to a general polynomial, which incorporates recoil/radiation reaction, in the emitted photons frequency, thereby connecting it with the Unruh effect. The presence of the Unruh effect is explored via a Rindler frame analysis, which also provided a path to directly measure the FDU temperature. This work not only explores channeling radiation, in a quantitative manner, but also sheds new light on it and the Unruh effect in a manner that is backed by the experimental evidence. Moreover, by analyzing the Rindler horizon thermodynamics, we were also able to confirm the Bekenstein-Hawking area-entropy law. In conclusion, our analysis indicates that in addition to measuring radiation reaction, the recent high-energy channeling radiation experiment has a significant statistical signal for the first observation of acceleration-induced thermality backed by an independent thermodynamic analysis of the Bekenstein-Hawking area-entropy law, thereby indicating the first observation of the Unruh effect in a nonanalog system.

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