

A Strained Crystal Is a Designer Spacetime

Brecht Corbeel

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Abstract. Strain enters the low-energy action of a Dirac or Weyl material exactly as a gauge or metric perturbation, so a machined strain profile is a designer spacetime for its electrons. The calling card is already on the bench: strained graphene nanobubbles realise pseudo-magnetic fields above 300 T [50], an order of magnitude past any laboratory magnet. We assemble the strain-to-geometry dictionary as a quantitative engineering table across named materials, then design a concrete horizon device — a near-critical type-II Weyl slab (the $\text{WTe}_2/\text{MoTe}_2$ family) whose bending-beam strain gradient sweeps the cone tilt W through the type-I/type-II boundary, so the line where $W = 1$ is an electronic horizon with analog surface gravity $\kappa_{\text{eff}} = v_F |\partial W / \partial x|$ and analog Hawking temperature $T_H^a = \hbar \kappa_{\text{eff}} / 2\pi k_B$. At a fabricable design point (0.3% strain over 100 nm) $T_H^a \simeq 0.1$ K, in the dilution-refrigerator range, spanning 0.04–0.4 K across the design envelope. The honesty is in two places. First, no experiment has yet measured the tilt’s strain derivative $\partial W / \partial u$ in any real type-II material [48]; every strain→tilt number here is computed (DFT/tight-binding), and measuring it is the device’s first milestone. Second, the kill condition: horizon physics survives only where the carrier stays ballistic across the horizon, $\mu > \mu_{\text{min}}(T_H^a) = ev_F / k_F k_B T_H^a$, and at the design point the near-critical type-II hosts that supply the tilt ($\mu \sim 10^3\text{--}10^4$ cm²/Vs) sit two orders of magnitude below threshold, while the crystals that clear it (Cd_3As_2 , $\mu > 10^7$; NbP , 5×10^6) are type-I and carry no tilt — so no material today clears both bars, and the program is a few-kelvin design in demonstrably cleaner type-II crystals. We give the signature menu at the design point — a shot-noise thermal factor at T_H^a , a two-terminal conductance peak of order $2e^2/h$ per mode, and scanning thermometry at the SQUID-on-tip floor — each a kinematic-analog quantity, in the sense the companion paper fixes [13]; the paper never claims gravity.

Keywords: pseudo-magnetic fields; strain engineering; Dirac and Weyl semimetals; type-II Weyl; tilted cones; analog horizons; elastic gauge fields; graphene.

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1 Strain already bends electronic spacetime by three hundred tesla

Strain enters the low-energy action of a Dirac or Weyl material as a gauge or metric perturbation of its electrons, so a machined strain profile is a designer spacetime for the carriers that live in it. The correspondence is exact at leading order in the $k \cdot p$ /tight-binding expansion — a smooth deformation shifts and tilts the Dirac cone in precisely the way a background vector potential or an off-diagonal metric component would [42, 78, 83] — and we grade it as exact only at that order: the full tight-binding treatment carries additional terms that do not couple as a standard gauge field and are of the same order of magnitude as the ones that do [19], so every analogical step in this paper is graded exact, structured, or speculative in the sentence where the step is made.

The calling card is already on the bench. Scanning tunnelling microscopy of strained graphene nanobubbles reads pseudo-magnetic fields in excess of 300 T [50], with

theory placing the peak curvature regions at 200–400 T — an order of magnitude past the ~ 85 T that the strongest sustained laboratory magnet reached at the time, and reached without a magnet at all. The record is not a fragile one-off: architected nanopillar substrates drive graphene to ~ 800 T at room temperature [33], and heteroepitaxial nanoprism arrays produce a uniform 41 T pseudo-field whose Landau levels are resolved by photoemission at room temperature on a wafer scale [62]. When the object of the analogy is the electron’s own dispersion, “geometry” is not a metaphor and not small: it is the largest effective field in condensed-matter physics.

This paper takes that fact seriously enough to build a device from it, and it inherits an honest amount to do so. The strain–gauge dictionary is the graphene community’s [2, 27, 42, 78, 83]; the 300 T measurement is Levy et al.’s [50]; the type-II Weyl phase, whose over-tilted cone is the geometry we exploit, is Soluyanov et al.’s [74]; and the idea that an over-tilted region reproduces a black-hole horizon has a fifteen-year literature, from Volovik’s ancestor claim [81] through the type-III/IV classification [63], the junction mappings [43], the lattice-

model series [46, 56, 59, 71], and the transport-signature work [18, 31], with the helium-3 program as its own older root [36, 38]. The analog-gravity program that licenses all of it is Unruh’s and its modern review is Barceló, Liberati and Visser’s [5, 79].

We declare the grade of every “spacetime,” “horizon,” and “Hawking” claim up front, because the honesty of the whole construction depends on it. Following the partition fixed by the companion paper [13], an analog observable is *kinematic* when it is a functional of the effective metric, the field’s quantum state, and the dispersion relation alone, and *dynamic* when its black-hole counterpart requires the Einstein field equations, their matter sources, or the area–entropy law. Our device engineers the electronic band-geometry, which is exactly the object the analogy is kinematic about — the dispersion and the null-cone structure it sets — and it is silent on everything the analogy is not: there is no Bekenstein–Hawking entropy here, no evaporation lifetime or Page curve, and the “back-reaction” of the electrons is the crystal’s elasticity and screening, not spacetime sourced by Einstein’s equations [13]. Every signature in Section 5 is kinematic in this precise sense, and we say so at each one. The paper never claims gravity.

Against that inherited backdrop, four things are ours, and a verified audit of eighteen horizon proposals spanning 1998–2023 found no prior paper combining even two of them. We (i) assemble the strain-to-geometry dictionary as a quantitative engineering table across named materials, every entry sourced (Section 2); (ii) design a concrete device — a flexure-strained near-critical type-II Weyl slab whose tilt gradient sweeps the cone through the type-I/type-II boundary — and compute its analog surface gravity κ_{eff} and analog Hawking temperature T_H^a from real material parameters (Section 3); (iii) derive the disorder kill condition, a mobility and mean-free-path threshold below which the horizon physics is buried, and state it as a number the field can aim at (Section 4); and (iv) compute the magnitudes of three experimental signatures at the design point, each graded for its analog-partition status (Section 5). The closest prior work is De Beule et al.’s heterostructure-interface horizon, which pairs an abstract-model T_H^a with a genuine conductance-peak signature but names no material and derives no disorder threshold [18]; Huang, Jin and Liu name a real compound ($\text{Zn}_2\text{In}_2\text{S}_5$) but compute neither a temperature nor a signature [35]; and Volovik’s ancestor gives only the symbolic relation $T_H^a = (\hbar/2\pi k_B) \partial v / \partial r$ with no material, no number, and no fabrication route [81]. The four-part package, anchored to one strained near-critical crystal, is unattempted.

Two admissions carry the paper’s weight, and both appear before any device is drawn. First, the tilt’s strain derivative $\partial W / \partial u_{ij}$ has never been measured: no experiment from 2016 to 2026 reports a tilt parameter under strain in any real type-II material, and the state of the art reaches only 0.34% operando

strain on $\text{Mo}_{0.91}\text{W}_{0.09}\text{Te}_2$ while reporting band and transport changes without a tilt number [48]. Every strain-to-tilt entry here is therefore computed (density-functional/tight-binding), marked as such, and measuring it directly is the device’s stated first milestone. Second, the kill condition bites today: the near-critical type-II hosts that supply the tilt carry mobilities two orders of magnitude below what a sub-kelvin horizon demands, while the ultrahigh-mobility crystals that would resolve it are type-I and carry no tilt to bend, so no material now clears both bars at once. The dictionary is real, the device geometry is fabricable, and the honest boundary is the subject of Section 4 rather than a footnote to it.

2 Strain enters the low-energy action as gauge and metric fields, entry by entry

The dictionary that turns a deformation into a field has two columns, one for the two-dimensional Dirac matter where it was written and one for the three-dimensional Weyl matter where this paper puts it to work. Both follow from the same move: expand the lattice Hamiltonian to leading order in a slowly varying strain field $u_{ij}(\mathbf{r})$ and read off how the deformation dresses the Dirac operator. We state the two columns as a single proposition, transcribing the coupling forms from the sources that derived them, and grade the correspondence exact at this order.

Proposition 1 (The strain-to-geometry dictionary; exact within $k \cdot p$ /tight-binding order, attributed). *For a two-dimensional Dirac material (graphene), a strain field couples to the low-energy electrons as a pseudo-gauge potential*

$$A_{\text{ps}} = \frac{\hbar\beta_G}{2ea} (u_{xx} - u_{yy}, -2u_{xy}), \quad \beta_G \approx 2-3, \quad (1)$$

with a the lattice constant and β_G the Grüneisen-class electron–phonon coupling, and its curl is a pseudo-magnetic field $B_{\text{ps}} = \nabla \times A_{\text{ps}}$ sourced by strain gradients; the measured anchor is the nanobubble field $B_{\text{ps}} > 300 \text{ T}$ [50]. For a three-dimensional Weyl semimetal, the same strain both displaces the nodes — a chiral (axial) gauge field $\mathbf{b}(\mathbf{r})$ — and tilts the cones, so that the low-energy Hamiltonian

$$H = v_F \boldsymbol{\sigma} \cdot (\mathbf{k} - \mathbf{b}(\mathbf{r})) + v_F W(\mathbf{r}) \cdot \mathbf{k} \not\ll \quad (2)$$

describes an electron in an effective metric of Painlevé–Gullstrand form wherever the tilt $W(\mathbf{r})$ varies, the type-II condition being $|W| > 1$ (Appendix A, conventions fixed). The engineering table gives, per material, the measured or computed $\partial W / \partial u_{ij}$, $\partial \mathbf{b} / \partial u_{ij}$, v_F , and achievable elastic strain, with sources.

Equation (1) is the closed form the graphene reviews settle on: a strain-tensor combination scaled by the dimensionless coupling $\beta_G = -\partial \ln t / \partial \ln a \approx 2$, with t the

nearest-neighbour hopping and $a = 1.42 \text{ \AA}$ [2, 11, 83]. The coupling is not perfectly pinned, and we carry the spread honestly: full tight-binding fits to uniaxial strain place the effective vector-potential coupling in the band $\beta_G \approx 2.6\text{--}4.0$ depending on strain direction and which hopping term is tracked [66], a quantity that must be kept separate from the *phonon* Grüneisen parameter $\gamma(E_{2g}) \approx 1.99$ measured from the Raman G-peak splitting under strain [58], with which it is frequently conflated. That the map (1) is exact only at leading order is itself a result: a systematic tight-binding derivation finds an additional strain-induced term that does not couple to the electrons as a standard gauge field and is of the same order of magnitude as the ones in Eq. (1) [19], which is why we grade the pseudo-gauge correspondence structured rather than exact once strain gradients are large.

The measured pseudo-field record is real, large, and increasingly designed. The graphene column is the one part of the dictionary with a decade of measurements behind it, and they run from single bubbles to wafers. Levy et al.’s nanobubbles on Pt(111) set the headline at $B_{ps} > 300 \text{ T}$ over lateral sizes of 4–10 nm [50]; architected nanopillar substrates push a strain-free monolayer to $\sim 800 \text{ T}$ at room temperature over a scalable area [33]; and heteroepitaxial SiC nanoprisms deliver a *uniform* 41 T whose pseudo-Landau levels are resolved by photoemission at room temperature [62]. The field can also be dialled: atomic-force-microscope-programmable nanobubbles span 0–125.7 T with independently tunable size and location [39], competing tip and gate forces on a suspended drumhead confine carriers into strain-defined quantum dots [45], an hBN moiré “magnifying glass” calibrates the local strain that produces the field [41], and periodic strain arrays generate a quantization distinct from the ordinary Landau ladder of a single smooth bump [4]. This is the empirical warrant for the whole paper: when the field acts on the electron’s own dispersion, engineered strain routinely reaches hundreds of tesla, so the question for a horizon device is not whether strain-geometry is strong enough but whether the *three-dimensional tilt* column of the dictionary can be filled with the same confidence — and it cannot yet, for the reason the next paragraph makes precise.

The Weyl column is where the honesty lives. The chiral gauge field $\mathbf{b}(\mathbf{r})$ in Eq. (2) has an explicit closed form on a tight-binding Weyl lattice, linear in the strain tensor through a Grüneisen coupling [14], and its spatial variation acts as an axial vector potential whose $n=0$ pseudo-Landau level reproduces the Fermi arc [25, 68]. The tilt term $v_F W(\mathbf{r})$ is the one we exploit, and it is the term the literature has never measured under strain. The type-II phase itself is on firm ground: it was predicted for WTe_2 [74] and MoTe_2 [77, 84] and observed by photoemission in MoTe_2 [21, 34, 40] and WTe_2

[85], with the tilt criterion — constant-energy surfaces opening from closed ellipsoids to hyperboloids once the cone velocity offset exceeds the cone slope — set out in the field review [3]. But the reading of an observed Fermi arc as automatic proof of the type-II state is contested: Bruno et al. showed that WTe_2 carries a large, *topologically trivial* surface Fermi arc that exists independently of the bulk Weyl points [8], so the arc’s presence neither confirms nor rules out the tilt, and we treat WTe_2 ’s near-critical status as computed rather than measured on that account too.

What has never been measured is the derivative that the device needs: $\partial W/\partial u_{ij}$, the rate at which strain changes the tilt. The most recent operando-strain photoemission study reaches 0.34% tensile strain on $\text{Mo}_{0.91}\text{W}_{0.09}\text{Te}_2$, reversible below 0.17%, and measures a $\sim 30\%$ increase in the density of states at E_F — but reports no tilt parameter and locates no type-I/type-II crossing, validating its bands only against *ab initio* calculation [48], and the group reviews confirm no such derivative has been compiled for any type-II material [37]. Every strain-to-tilt number in this paper is therefore computed, and we mark it so. We anchor the computation on the density-functional strain sensitivities of the type-II literature — a 0.5% compression that drives the WTe_2 Weyl-point separation from $\sim 0.7\%$ to $\sim 2\%$ of the reciprocal lattice vector [74], and the DFT-predicted strain-driven type-II/type-I transitions of MoTe_2 [77, 84] — together with a Grüneisen estimate $\partial W/\partial u \sim W \beta_G$ using $\beta_G \approx 2\text{--}3$. Reading a boundary crossing $\Delta W \sim 0.1\text{--}0.3$ over $\sim 0.5\text{--}2\%$ strain, both routes land in the band $\partial W/\partial u \sim 5\text{--}30$, central value 15; this is a genuinely novel contribution rather than a restatement of measured data, and Appendix B gives its provenance entry by entry. That 8-Pmmn borophene supplies a tilted-Dirac cross-check model with a closed-form strain coupling of the same family [89] is reassuring for the form, not for the number.

The dictionary as an engineering table. Table 1 collects the shortlist: the two untilted anchors that bracket the problem (graphene, the pseudo-field champion and mobility benchmark; Cd_3As_2 and NbP, the ultrahigh-mobility Dirac and Weyl crystals), the near-critical type-II hosts that carry the tilt (the $\text{WTe}_2/\text{MoTe}_2$ family), and the borophene-class model. Figure 1 renders the two-column logic as a plate. The reason no single row is both tilted and clean is the whole content of Section 4; the table simply makes the trade-off legible, with every measured number carrying a primary source and every computed tilt sensitivity flagged as computed. The Fermi velocity entered for the type-II hosts is the design host value used throughout the device arithmetic ($v_F \simeq 2 \times 10^5 \text{ m/s}$, the reduced along-tilt velocity of a near-critical band; Appendix C), not a separately measured constant, and we mark it as such.

The table is the paper’s ledger. It says, in one place,

Table 1: The material shortlist for a strain-engineered electronic horizon. Measured quantities carry primary sources; the tilt sensitivity $\partial W/\partial u$ is computed (density-functional/tight-binding), never measured in any type-II material to date [48], with provenance in Appendix B and the propagated range in Appendix C. “Host” v_F is the design host value of the near-critical along-tilt band used in the device arithmetic, not a separate measurement.

Material	Type	v_F (m/s)	Measured (cm^2/Vs)	μ	Rest W	tilt	$\partial W/\partial u$ (computed)	Achievable strain	Sources
Graphene	I (2D Dirac, untitled)	$\sim 10^6$	$> 2 \times 10^5$ (susp.)	0		— (untitled)	— (untitled)	$\sim 6\%$ uniform; $\sim 25\%$ local	7, 10, 11, 49
WTe ₂	II (near-critical) (host)	$\sim 2 \times 10^5$	$\sim 10^3\text{--}10^4$	$\gtrsim 1$		5–30 (central 15)	$\sim 0.23\%$ (piezo)		1, 32, 74
MoTe ₂	/ II (near-critical) (host)	$\sim 2 \times 10^5$	$\sim 10^4$	$\gtrsim 1$		5–30 (central 15)	0.34% operando		48, 77, 84
MoWTe ₂	I (Dirac, untitled)	$\sim 1.5 \times 10^6$	$> 10^7$	0		— (untitled)	$\sim 3\text{--}8\%$ (torsion, modeled)		51, 68
Cd ₃ As ₂	I (Weyl, untitled)	—	5×10^6	< 1		— (untitled)	—		24, 73
NbP	I (Weyl, untitled)	—	5×10^6	< 1		— (untitled)	—		24, 73
8-Pmmn borophene	tilted Dirac (model)	model	—	< 1		coupling form (model)	—		89

that the geometry knob is strong (the graphene row), that the tilt exists in real crystals (the type-II rows), that the tilt sensitivity is computed and not yet measured (every dash and every 5–30), and that the crystals which would let a horizon be seen are the ones without a tilt to bend (the Cd₃As₂ and NbP rows). The device of Section 3 is built from the type-II rows; the verdict of Section 4 is the arithmetic of why those rows, today, sit on the wrong side of the disorder threshold.

3 A bent slab of near-critical type-II Weyl is an electronic horizon

The device is the dictionary’s Weyl column turned into a fabrication drawing. Take a slab whose cones are just over-tilted at rest, bend it so that a strain gradient runs along a short channel, and the tilt falls smoothly through the critical value; the line where it crosses is a place past which one branch of carriers can only move one way, which is what an electronic horizon is. The construction is ours, and we state it as a proposition with its surface gravity fixed to the normalization derived in Appendix A.

Proposition 2 (The horizon device; ours). *Take a suspended or flexure-strained slab of a near-critical type-II candidate — a material from Table 1 with rest tilt near $|W| = 1$, the WTe₂/MoTe₂ family leading — and deliver, by bending-beam flexure, a uniaxial strain gradient that sweeps W through 1 across a channel of length L . The crossing line is an electronic horizon with analog surface gravity*

$$\kappa_{\text{eff}} = v_F \left| \frac{\partial W}{\partial x} \right|_{W=1}, \quad (3)$$

the $O(1)$ normalization fixed by the analog surface-gravity

definition in Appendix A, and analog Hawking temperature

$$T_H^a = \frac{\hbar \kappa_{\text{eff}}}{2\pi k_B}. \quad (4)$$

With 0.1–1% strain delivered over 100 nm–1 μm and the shortlist’s computed $\partial W/\partial u$, Eqs. (3)–(4) place T_H^a in the 0.1–10 K range (numerical values from the seeded computation of Appendix C, plotted in Fig. 2).

We grade the horizon claim exact on the kinematic side and silent on the dynamic side, in the sense fixed in Section 1 [13]: Eq. (2) with a varying tilt is, at leading order, an electron propagating on a Painlevé–Gullstrand effective metric whose g_{0i} components are the tilt [63, 81], and the null-cone tipping at $W = 1$ is the same geometry Unruh’s sonic horizon and its condensed-matter descendants realize [5, 79]. The surface-gravity form $\kappa_{\text{eff}} = v_F |\partial W/\partial x|$ is Volovik’s tilt-gradient temperature [81] and coincides with De Beule et al.’s $T_H^a = \hbar v_F \alpha / 2\pi k_B$ once their tilt slope α is identified with $\partial W/\partial x$ at the crossing [18]; Appendix A shows the factor of order unity between the raw gradient and the properly defined analog surface gravity, so that no hidden normalization inflates T_H^a .

Why the tilt must be near-critical, and why that forces the material choice. The horizon needs W to pass through 1 inside the achievable strain window. If the rest tilt is far below 1, no laboratory strain reaches the crossing — the elastic limit runs out first; if it is far above 1, the crossing sits at a strain the crystal cannot survive. Only a near-critical host, $|W|$ already close to 1 at rest, brings the crossing within a fraction of a percent of strain. That is exactly the WTe₂/MoTe₂ family: type-II by prediction and photoemission [21, 34, 74, 77], sitting so close to the boundary that density-functional theory finds strain can drive them back across it [77, 84].

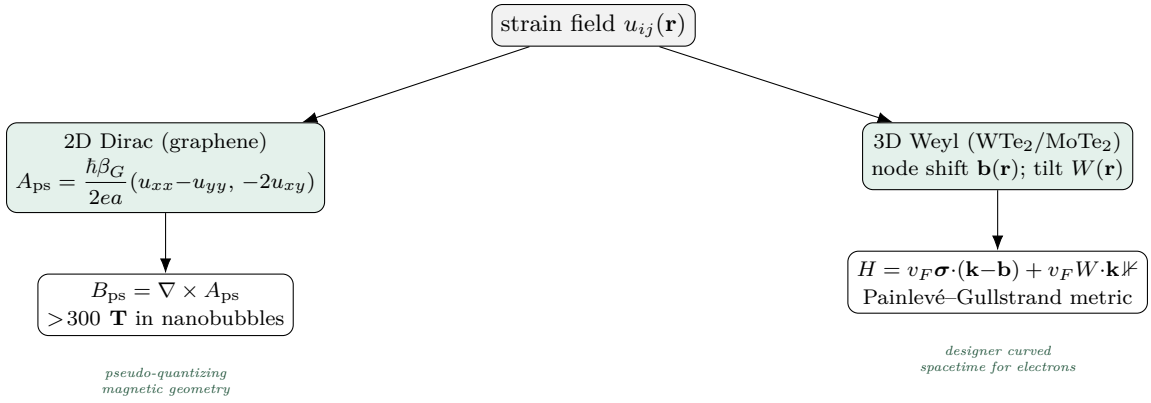


Figure 1: The strain-to-geometry dictionary (Proposition 1). A single machined strain field enters the low-energy action of a Dirac or Weyl material as a gauge or metric perturbation. In 2D Dirac matter it is a pseudo-gauge potential A_{ps} whose curl is a pseudo-magnetic field: strained graphene nanobubbles already realise $B_{\text{ps}} > 300$ T [50], an order of magnitude beyond any laboratory magnet. In a 3D Weyl semimetal the same strain both displaces the nodes — a chiral (axial) gauge field $\mathbf{b}(\mathbf{r})$ [14] — and tilts the cones, so that a spatially varying tilt $W(\mathbf{r})$ enters the Hamiltonian exactly as the off-diagonal components of a Painlevé–Gullstrand metric (Appendix A). The engineering table of Section 2 makes each entry quantitative, material by material; every strain→tilt figure is computed, not measured [48, 74].

The ultrahigh-mobility crystals of Table 1 are disqualified here not on cleanliness but on geometry: Cd_3As_2 and NbP are type-I, their cones untilted, so there is no $W \rightarrow 1$ crossing to engineer at all [51, 73], and suspended graphene, for all its pseudo-field prowess, is a two-dimensional Dirac system with no cone tilt to bend. We adopt $\text{Mo}_{0.91}\text{W}_{0.09}\text{Te}_2$ as the concrete lead, because it is the one near-critical type-II crystal on which calibrated operando strain has already been delivered and its band response measured [48], even though that experiment stopped short of a tilt parameter. We defend the choice with the caveat attached: WTe_2 ’s type-II status was challenged on the grounds that its observed Fermi arc is topologically trivial and tilt-independent [8], so “near critical” is, for now, a computed property of these hosts and the first thing an experiment must confirm.

The fabrication route is assembled from demonstrated parts. Nothing in the device asks for a process that has not been run. The strain is delivered by a piezoelectric bending-beam cell of the Hicks–Mackenzie class, which reaches continuously tunable, reversible strain up to $\sim 0.23\%$ at cryogenic temperature with good homogeneity on a high-aspect-ratio sample [32]; the operando-strain photoemission cell that took $\text{Mo}_{0.91}\text{W}_{0.09}\text{Te}_2$ to 0.34% shows the type-II family tolerates comparable strain reversibly [48], and free-standing membranes of the two-dimensional analogue sustain several percent uniform elastic strain before fracture [10, 49]. The slab itself follows the suspended-membrane route of the graphene electromechanical-resonator lineage — exfoliation over a trench, clamped at the contacts, released to bend [9]; a boron-nitride substrate or encapsulation supplies the atomically smooth, low-disorder support that

raised graphene mobility by nearly an order of magnitude [17, 55], which Section 4 will show is the binding constraint. Ohmic source and drain contacts to a transition-metal ditelluride are available by laser phase-patterning of a metallic $1T'$ region into the semiconducting host, which raised MoTe_2 transistor mobility about 50-fold while holding an on/off ratio of 10^6 [12]. Over the channel sits a scanning SQUID-on-tip thermometer of sub-50 nm diameter for the local-dissipation readout of Section 5 [29]. The layout — cross-section, strain map, contacts, and the engineered $W(x)$ profile — is drawn in Fig. 2.

The contact and readout layout follows the measurement. The channel is two-terminal by design, because the cheapest decisive signature is a two-terminal conductance peak (Section 5): source and drain bracket the crossing, the number of transverse modes is set by the channel width, and the phase-patterned $1T'$ contacts keep the injection ohmic [12]. The strain map is the other half of the drawing and is fixed by beam mechanics: bending a slab of thickness d to radius of curvature R delivers a surface strain $u \simeq d/2R$ that varies along the channel, so the tilt gradient $\partial W/\partial x$ follows directly from the beam geometry and the computed $\partial W/\partial u$, and is uniform to the extent the curvature is. The choice of local bending over the sample-wide routes that blanket tens of micrometres by heterostructure lattice mismatch [52] is deliberate: a horizon needs a controlled gradient across a short channel, not a uniform field over a large area, and the gradient is exactly the quantity Eq. (3) turns into surface gravity. The scanning thermometer sits above the channel for the local-dissipation image, so the same device serves all three readouts of Section 5

without refabrication [29].

The design point, and the envelope around it.

Fixing the arithmetic to one fabricable operating point makes the claim checkable. We take a host along-tilt Fermi velocity $v_F = 2 \times 10^5$ m/s and a strain of 0.3% delivered over $L = 100$ nm; with the central computed sensitivity $\partial W/\partial u = 15$ this gives a tilt gradient $\partial W/\partial x = 0.45 \mu\text{m}^{-1}$, hence from Eq. (3) a surface gravity $\kappa_{\text{eff}} = 9.0 \times 10^{10} \text{ s}^{-1}$ and from Eq. (4) an analog Hawking temperature $T_H^a = 0.11$ K (Appendix C; Fig. 2). This is squarely in the dilution-refrigerator range, which is the point: the temperature is low, but not fantastically low, and every input is a number the shortlist supplies. The envelope is wide enough to matter and bounded enough to be honest. Sweeping strain over 0.1–0.5%, channel length over 50 nm–1 μm , and the computed sensitivity over its 5–30 band, T_H^a spans 3.6 mK–0.36 K (Appendix C). Holding the design geometry and sweeping only the sensitivity across its conservative, central, and optimistic values gives $T_H^a = 0.037, 0.109$, and 0.219 K — a factor of six that propagates directly from the one quantity the paper admits is computed rather than measured, $\partial W/\partial u$, which is why measuring it is the program’s first milestone (Section 6).

The arithmetic that produces a tenth of a kelvin is the easy half of the design. It says the horizon is warm enough to reach in a standard cryostat and cold enough that its signatures live below the ordinary thermal floor. What it does not say — what no earlier proposal in the horizon literature said, because none carried a material and a mobility — is whether a carrier in this particular slab survives long enough to feel the horizon before disorder scatters it. A surface gravity of $9 \times 10^{10} \text{ s}^{-1}$ sets a Hawking time $\hbar/k_B T_H^a \approx 70$ ps and, through the host velocity, a horizon length of order ten micrometres; whether the mean free path clears that length is the difference between a designer spacetime and a warm, disordered resistor. That comparison is Section 4, and it is where the design meets its verdict.

4 Disorder buries the horizon before any present material can show it

A horizon is a coherent kinematic structure, and disorder is its solvent. For the tipping of the null cone to be a physical fact for a carrier rather than a line on a band diagram, the carrier must propagate ballistically across the crossing for at least as long as the horizon takes to act. That single requirement, turned into an inequality on mobility, is the paper’s engineering teeth, and it is where every prior proposal stopped. We state it, evaluate it against measured mobilities, and give the verdict without softening it.

Proposition 3 (The kill condition; ours). *Horizon physics requires the Hawking time $\hbar/k_B T_H^a$ to be shorter than the transport scattering time $\tau = \mu m^*/e$, equivalently the mean free path $\ell = v_F \tau$ to exceed the horizon length $\ell_H \equiv \hbar v_F/(k_B T_H^a)$. With the Weyl transport mass $m^* = \hbar k_F/v_F$ this is a threshold on mobility,*

$$\mu > \mu_{\min}(T_H^a) = \frac{ev_F}{k_F k_B T_H^a}, \quad (5)$$

a curve $\mu_{\min}(T_H^a)$ per material. Against measured mobilities the verdict follows: which shortlist materials clear the bar today, which need cleaner growth, and by how much.

The derivation is one line and it fixes the physics of the bound. The horizon acts over the Hawking time $t_H = \hbar/k_B T_H^a$; a carrier resolves it only if it survives ballistically for longer, $\tau > t_H$. Writing the transport time as $\tau = \mu m^*/e$ and the linear-band transport mass as $m^* = \hbar k_F/v_F$ turns $\tau > \hbar/k_B T_H^a$ directly into Eq. (5), while multiplying through by v_F restates it as the geometric condition $\ell > \ell_H$: the mean free path must be longer than the region over which the metric bends. The two forms are the same statement, and both scale as $1/T_H^a$ — a colder horizon is a longer, more fragile one. This is the threshold no paper in the horizon literature derived: De Beule et al. came closest, citing TaAs’s measured mean free path of $\sim 5.2 \mu\text{m}$ once as a plausibility remark, but never turned it into a critical value [18], and the group review states outright that a complete scattering theory in the presence of these pseudo-fields is still absent [37].

Which lifetime enters, and at what density, decides how the bound is read. Equation (5) is written in the transport scattering time, and that choice is deliberate: the horizon is a real-space structure a carrier must cross ballistically, so it is the transport mean free path — not the typically much shorter quantum lifetime that sets the decay of quantum oscillations — that must clear ℓ_H [16]. In the ultrahigh-mobility crystals the two differ by orders of magnitude, the transport lifetime exceeding the quantum lifetime by $\sim 10^4$ in Cd_3As_2 [51], which is why those crystals clear the bar so comfortably despite non-negligible single-particle scattering. The threshold also turns on carrier density through k_F : since $\mu_{\min} \propto k_F^{-1}$, running at higher density relaxes the mobility demand, but only by pushing E_F off the node and diluting the tilt physics the device depends on, so we fix the representative $n \sim 10^{18} \text{ cm}^{-3}$ used throughout rather than buy margin that way. One further caveat sharpens the target rather than softening it: near the node the carriers can enter a hydrodynamic regime where momentum-conserving collisions dominate [53], and a viscous electron fluid is not a ballistic carrier crossing a horizon, so clearing Eq. (5) is the necessary condition, not the sufficient one.

At the design point the number is severe. Evaluating Eq. (5) at $T_H^a = 0.11$ K, with k_F set by a representative

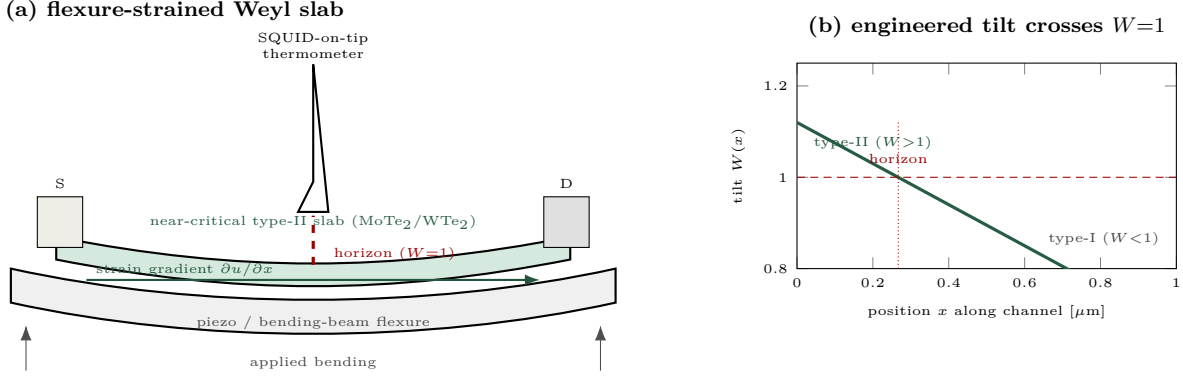


Figure 2: The horizon device (Proposition 2). **(a)** A slab of a near-critical type-II Weyl semimetal (accent) is bonded to a bending-beam flexure; the applied curvature delivers a uniaxial strain gradient $\partial u / \partial x$ along a channel of length $L \sim 100$ nm between source and drain, with a scanning SQUID-on-tip thermometer over the channel. **(b)** Through the computed tilt sensitivity $\partial W / \partial u$ (Section 2), the gradient sweeps the tilt $W(x)$ from over-tilted (type-II, $W > 1$) to under-tilted (type-I, $W < 1$); the line where $W = 1$ is an electronic horizon with analog surface gravity $\kappa_{\text{eff}} = v_F |\partial W / \partial x|$ and analog Hawking temperature $T_H^a = \hbar \kappa_{\text{eff}} / 2\pi k_B$. At the design point (strain 0.3% over 100 nm, central sensitivity) $T_H^a \simeq 0.11$ K; the full envelope and the normalization of κ_{eff} are in Appendices A and C.

carrier density $n \sim 10^{18} \text{ cm}^{-3}$, gives a required mobility $\mu_{\text{min}} \approx 7 \times 10^5 \text{ cm}^2/\text{Vs}$ and a horizon length $\ell_H \approx 14 \mu\text{m}$ (Appendix C; Fig. 3). A carrier must travel fourteen micrometres without scattering to register a horizon a tenth of a kelvin warm. Table 2 sets that demand against what the shortlist crystals actually deliver.

The verdict is plain, and we state it plainly: today no material offers both a near-critical type-II tilt and the mobility to resolve a sub-kelvin horizon. The three near-critical hosts that carry the tilt — WTe₂, MoTe₂, and the strain-tunable Mo_{0.91}W_{0.09}Te₂ — sit at mobilities of order 10^3 – $10^4 \text{ cm}^2/\text{Vs}$ [1, 48], roughly two orders of magnitude below their own thresholds; the crossing they would show is buried under scattering long before a carrier reaches it. The crystals that clear the mobility bar with room to spare, Cd₃As₂ at $\mu > 10^7$ [51] and NbP at 5×10^6 [73], are type-I: their cones are not tilted, so there is no $W = 1$ crossing to bend into a horizon. Suspended graphene, the mobility champion of the two-dimensional world at $\sim 2 \times 10^5 \text{ cm}^2/\text{Vs}$ [7], fails on both counts, being both short of its threshold and untilted. The gap between the tilted column and the clean column is the whole result: the two properties the device needs live in different crystals.

The bound is not a wall, and its slope is the way through. Because μ_{min} falls as $1/T_H^a$, a warmer horizon is a cheaper one to resolve: pushing the design toward $T_H^a \sim 10$ K — by a larger tilt sensitivity, a shorter channel, or a stiffer strain gradient — drops the required mobility to $\sim 10^4 \text{ cm}^2/\text{Vs}$ (Fig. 3; Appendix C), which is within reach of the type-II hosts’ present growth rather than two orders beyond it. The kill condition therefore does not close the program; it aims it. It says that the first demonstration will not be a sub-kelvin horizon in today’s WTe₂, but a few-kelvin horizon in a demonstra-

bly cleaner type-II crystal — and it says, as a number, exactly how much cleaner: the mobility of a candidate host must clear Eq. (5) at whatever T_H^a its strain envelope allows. That is the target Section 6 hands to crystal growth, and the verdict here is the reason the target is stated in mobility rather than in hope.

5 Three measurements can see the horizon, and one of them is cheap

Suppose the kill condition of Section 4 is met, in a cleaner crystal or a warmer design. Distinguishing the crossing from an ordinary resistive smear then falls to three observables, at three very different costs, and we compute the magnitude of each at the design point rather than assert that it exists. Each is a kinematic-analog quantity in the exact sense of Section 1, and we label it as such where it is stated.

Proposition 4 (Signatures; ours as computation, partition-graded). *At the design point the horizon produces three measurable signatures, each kinematic. (i) Two-terminal shot noise across the crossing carries a thermal factor set by T_H^a rather than the bath, so the excess noise temperature is of order T_H^a and is resolved for a bath temperature $T_{\text{bath}} \lesssim T_H^a$. (ii) The two-terminal differential conductance shows a peak of order $2e^2/h$ per transverse mode at the horizon, the Hawking-like transmission $T_{1,2} = 1/(1 + e^{\pm 2\pi\omega/\alpha})$ integrated over modes, whose position tracks the tilt slope $\alpha = \partial W / \partial x$ [18]. (iii) Scanning thermometry images the horizon dissipation, whose power at the design point is set against the SQUID-on-tip readout floor. Magnitudes are computed in Appendix C*

Table 2: The verdict. Required mobility $\mu_{\min}$ from Eq. (5) at the $T_H^a = 0.11$ K design point, each evaluated with the material’s own v_F and a representative $n \sim 10^{18} \text{ cm}^{-3}$ (Appendix C), against measured mobilities: WTe₂ from magnetotransport of the family [1], the strain-tunable Mo_{0.91}W_{0.09}Te₂ [48], Cd₃As₂ [51], NbP [73], and suspended graphene [7]. No material clears both the tilt requirement and the mobility bar.

Material	Near-critical type-II tilt?	Measured (cm^2/Vs)	μ	Required (cm^2/Vs)	$\mu_{\min}$	Verdict
WTe ₂	yes	5×10^3		6.9×10^5		fails mobility bar ($\sim 100\times$ short)
MoTe ₂ (Td)	yes	1×10^4		8.6×10^5		fails mobility bar ($\sim 100\times$ short)
Mo _{0.91} W _{0.09} Te ₂	yes	8×10^3		6.9×10^5		fails mobility bar ($\sim 100\times$ short)
Cd ₃ As ₂	no (type-I)	$> 10^7$		5.1×10^6		clears mobility; no tilt to bend
NbP	no (type-I)	5×10^6		1.0×10^6		clears mobility; no tilt to bend
Graphene (suspended)	no (untilted)	2×10^5		5.9×10^6		fails; no tilt to bend

and plotted in Fig. 4; all three sit on the kinematic side of the analog partition [13].

Shot noise is the thermometer of the emission, and it is the expensive one. Charge quantization makes current noisy, and the noise carries information the average conductance does not — the potential profile, the relevant energy scales, and, here, the temperature of the emission across the crossing [6]. The horizon adds an excess noise temperature of order T_H^a : a thermal factor set by the analog Hawking temperature rather than by the bath, which is the transport reading of Hawking-like emission across an electronic horizon. This is the most direct of the three signatures and the most demanding, because the excess is visible only where it rises above the ordinary Johnson–Nyquist floor, that is for $T_{\text{bath}} \lesssim T_H^a \simeq 0.11$ K (Fig. 4; Appendix C) — a dilution-refrigerator measurement, not a cryostat one. The instrumentation exists: microwave Johnson-noise thermometry resolves the electron temperature to $4 \text{ mK}/\sqrt{\text{Hz}}$ over an 80 MHz bandwidth [23], and the same technique read the electron system of graphene directly to expose the Dirac fluid [15], so a 0.1 K excess is well inside reach once the bath is cold enough. The signature has a distinctive shape: as the bath is cooled below T_H^a , the measured noise temperature stops tracking the bath and saturates on a plateau near T_H^a (Fig. 4), and the optimistic sensitivity $T_H^a \simeq 0.22$ K widens that window to temperatures a dilution refrigerator holds comfortably. We grade this signature kinematic: Hawking-type emission and its correlation functions are functionals of the effective metric, the field’s state, and the dispersion alone, and are proved to sit on the kinematic side of the partition [13].

The conductance peak needs no cold bath, and it is the cheap one. De Beule et al. showed that a slowly varying tilt horizon transmits carriers through Hawking-like thermal channels $T_{1,2} = 1/(1 + e^{\pm 2\pi\omega/\alpha})$,

and that the effect surfaces as a peak in the two-terminal differential conductance whose position tracks the tilt slope [18]. Summed over the transverse modes of the channel, the peak is of order $2e^2/h$ per mode — about $77 \mu\text{S}$ at the design geometry (Appendix C; Fig. 4) — a conductance-scale feature that a standard lock-in measures. Its decisive practical virtue is that, unlike the shot noise, it does not require cooling below T_H^a : the peak is a property of the transmission through the tipped cone, present as soon as the horizon exists, so it is the signature that a first, warmer device should target. We grade it kinematic on the same footing — it is a functional of the dispersion and the scattering geometry — while noting that its *interpretation* as “Hawking-like” is what the partition licenses, not the mere existence of a conductance peak, which disorder can also produce; the discriminating content is the peak’s tracking of the tilt slope, which Section 6 makes the second milestone. De Beule et al. give the position explicitly — for a slowly varying horizon the peak sits at $E_F/\omega_c \approx 0.7 \alpha/2k_0$, saturating at $E_F/\omega_c \approx 0.6$ for a sharp one [18] — so a peak that shifts with the applied strain gradient in the way the tilt slope $\alpha = \partial W/\partial x$ predicts is exactly the content a disorder-induced peak cannot reproduce.

Scanning thermometry is the picture, and it sits exactly at the floor. A SQUID-on-tip thermometer of sub-50 nm diameter and sub-microkelvin sensitivity images where power is dissipated [29, 30], which is what a horizon does at the crossing. The computed horizon dissipation at the design point is $\sim 2 \times 10^{-14} \text{ W}$ (Appendix C), against the instrument’s minimum detectable continuous power of $40 \text{ fW} = 4 \times 10^{-14} \text{ W}$ — the Landauer readout floor [29]. The signal is therefore about half the floor: the hardest of the three, feasible only with the design pushed toward larger dissipation (a warmer, more strongly bent horizon) or with a lower-noise probe. We report the ratio rather than round it away, because it is the honest statement of where the imaging signature

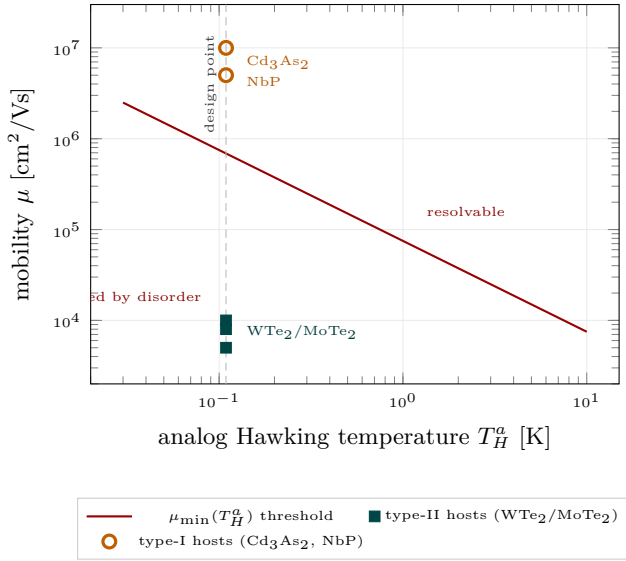


Figure 3: The kill condition, and why no material clears it yet (Proposition 3). Horizon physics survives only where the carrier stays ballistic across the horizon, $\mu > \mu_{\min}(T_H^a) = ev_F/(k_F k_B T_H^a)$ (red); the region above the curve is resolvable. At the device design point $T_H^a \simeq 0.11$ K (dashed) the near-critical type-II hosts that supply the tilt — WTe_2 , MoTe_2 and the strain-tunable $\text{Mo}_{0.91}\text{W}_{0.09}\text{Te}_2$ of Lanz et al. [48], at $\mu \sim 10^3\text{--}10^4$ cm^2/Vs — sit two orders of magnitude *below* threshold (filled), while the crystals that clear it, Cd_3As_2 ($\mu > 10^7$; 51) and NbP (5×10^6 ; 73), are type-I and carry no tilt to bend (open). The bar relaxes as T_H^a rises — a few-kelvin design brings the threshold down to $\sim 10^4$ cm^2/Vs , within reach of cleaner type-II growth — which is the program of Section 6. Values from the seeded computation of Appendix C.

stands. The precedent is demonstrated: the same probe class has imaged resonant dissipation from individual atomic defects in graphene [30], and Johnson-noise thermometry supplies a complementary, contactless readout of the electron temperature at 4 mK/ $\sqrt{\text{Hz}}$ [23] should the scanning route fall short. It, too, is kinematic: it images the effective geometry’s own dissipation structure, not spacetime’s response to radiated energy.

The three together are the argument. No single signature is decisive alone — a conductance peak can be disorder, an excess noise temperature can be stray heating, a dissipation hot-spot can be a contact — but the three carry different systematics, and their mutual agreement is the evidence. The conductance peak fixes the horizon’s position and its tracking of the tilt slope; the shot noise fixes the emission temperature and tests it against the computed T_H^a ; the thermometry fixes where the power is dissipated in real space. A device in which the peak sits where the thermometer sees the dissipation, at a strain the design predicts, while the noise saturates

at the T_H^a that the same strain gradient sets through Eq. (4), satisfies every kinematic test at once. That joint measurement, not any one curve, is what Section 6 proposes.

What “kinematic” buys, and what it forbids. All three signatures are kinematic, and the partition is not a hedge but a specification of exactly what the device does and does not test [13]. It buys the emission spectrum and its correlations, the transmission through the tipped cone, and the dissipation image — each a functional of the effective metric of Eq. (2), the carriers’ quantum state, and the dispersion, and each therefore a genuine analog of the corresponding black-hole quantity. It forbids everything the black-hole versions get from dynamics: there is no Bekenstein–Hawking entropy to measure here, because that quantity descends from the area law and the field equations; there is no evaporation lifetime or Page curve, because those are dynamical histories; and the device’s “back-reaction” is the response of the crystal — its elasticity, its screening, its own electrodynamics — not spacetime evolving under Einstein’s equations [13]. A measured agreement between the noise temperature and T_H^a would confirm the emission is thermal at the analog Hawking temperature; it would say nothing about quantum gravity, and we do not let the sentence drift toward implying otherwise. Other designed-geometry platforms have already measured Hawking-like emission on the kinematic side — sonic horizons in a Bose–Einstein condensate [47], an optical-fibre horizon [67], and a kelvin-scale polariton horizon [61] — and the solid-state device proposed here would join them as a transport-readable member of the same kinematic family [5], not as a claim about the other side of the partition.

6 What else the dictionary buys, and the two measurements that decide

The horizon is one entry in a longer table. Once strain is admitted as a geometry knob, the same dictionary designs other curved-space effects for electrons, some already measured, some only sketched, and the honest thing is to grade them and fence the speculative ones off from the device. We collect them as a proposition and keep the grades in the sentences.

Proposition 5 (Beyond the horizon; structured/speculative, graded). *The strain-to-geometry dictionary supports, beyond the horizon device: engineered “gravitational lensing” channels — strain lenses that deflect and split carriers by valley (structured, prior art cited); pseudo-Landau-level flat bands as designed quantizations (structured, with measured anchors); and a speculative row, analog cosmology in time-modulated*

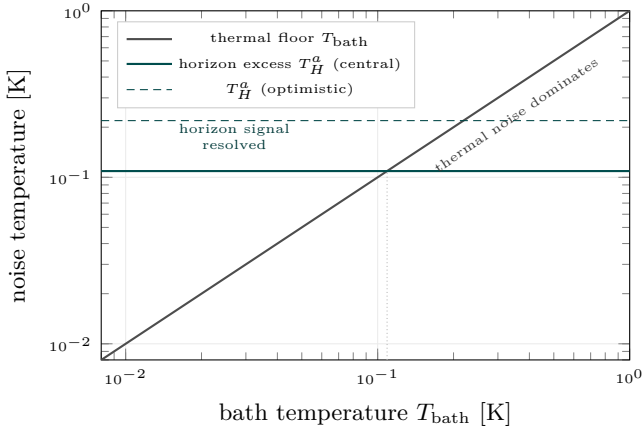


Figure 4: The signature window (Proposition 4). The horizon adds an excess shot-noise temperature of order T_H^a (teal), a thermal factor set by the analog Hawking temperature rather than the bath — the transport reading of Hawking-like emission across the electronic horizon. It is resolved only where it rises above the ordinary thermal floor T_{bath} (grey diagonal), i.e. below $T_{\text{bath}} \simeq 0.1\text{--}0.2$ K, placing the measurement squarely in the dilution-refrigerator range. The companion observable is the two-terminal differential-conductance peak of de Beule et al. [18], of order $2e^2/h$ per transverse mode ($\sim 77 \mu\text{S}$), which does not require sub- T_H^a cooling; scanning thermometry is the hardest of the three, since the $\sim 2 \times 10^{-14}$ W horizon dissipation sits at the SQUID-on-tip readout floor [29]. All magnitudes are kinematic-analog quantities (the sense fixed in Section 1 following 13); values from the seeded computation of Appendix C.

strain, which we fence and label. The program that would establish the horizon device itself is two measurements (structured).

Strain lenses are the lensing analog, and they already work as valley optics. A smooth strain profile bends carrier trajectories the way a gravitational potential bends light, and the effect is real enough to be a device: graphene nanobubbles act as valley filters and beamsplitters, sending K and K' electrons along different real-space paths with no external magnetic field [72]. In the tilted-Weyl setting the same idea is the “black-hole mirage” — tilt profiles that guide electron flow along paths analogous to gravitational lensing — and its authors are careful, showing by comparison to lattice simulation that the gravity analogy breaks down close to the analog horizon [31], which is exactly the grading discipline this paper adopts. A 2026 result reaches the same emergent-metric lensing from a different knob entirely, the Berry phase of a spatially varying spin texture, with deflection angles estimated at the milliradian scale [65]. We grade this row structured: the lensing is a genuine consequence of the effective geometry, the prior art is

cited, and the analogy is kinematic and bounded, not a claim about spacetime.

Flat bands are the designed quantization, and here the anchors are measured. When the pseudo-field is periodic rather than smooth, its pseudo-Landau levels flatten the bands by design. Buckled monolayer graphene forms a periodically strained superlattice whose flat bands host correlated states, an alternative to twisted bilayers requiring only one layer [54], and the theory companion shows how the strain geometry sets the flat-band width [57]. The measured anchors are strong: room-temperature pseudo-Landau levels at a uniform 41 T [62], a superlattice quantization distinct from the ordinary Landau ladder [4], and sample-wide pseudo-fields over tens of micrometres from a heterostructure lattice mismatch [52]. That a strain-driven topological-phase transition can be engineered into a piezo-topological transistor [26] rounds out the point: the dictionary’s non-horizon rows are, in several cases, already laboratory devices. We grade this row structured, with the quantization measured and the “designed spacetime” framing kept as the organizing analogy rather than a load-bearing claim.

The speculative row is analog cosmology, and we fence it. If static strain is a static metric, time-modulated strain is a time-dependent one, and a driven metric is the kinematic setting for particle production and analog cosmology. There is a real seed: an oscillating nano-deformation of a graphene bubble drives a pseudo-electric field and pumps a valley current at zero bias, in the sub-terahertz range [28], though its authors frame it as an oscillating deformation rather than a Floquet drive, and we keep that distinction. Engineered magnetoelectric media offer yet another designed-horizon platform in theory [20]. We grade this row speculative and fence it explicitly: no number in this paper depends on it, and the analog-cosmology reading of time-modulated strain is a direction we name, not a result we claim.

The program is two measurements, in order. The device’s fate turns on two experiments, and their order matters. The first is to measure the one quantity the paper admits is computed: $\partial W/\partial u_{ij}$, the tilt’s strain derivative. The state of the art applied calibrated operando strain to 0.34% on $\text{Mo}_{0.91}\text{W}_{0.09}\text{Te}_2$ and measured the band response, but stopped short of a tilt parameter and of locating the type-I/type-II crossing [48]; extracting $\partial W/\partial u$ and the crossing from an operando-strain photoemission sweep is the device’s first milestone, and it would replace every computed entry in Table 1 with a measured one, collapsing the factor-of-six spread that the 5–30 sensitivity band puts on T_H^a . The second is the conductance-peak measurement: the $\sim 2e^2/h$ -per-mode peak of Section 5, tracking the tilt slope [18],

measured in a near-critical type-II crystal grown clean enough to clear Eq. (5). Because the mobility bar falls as $1/T_H^a$, the honest first target is not a sub-kelvin horizon in today's crystals but a few-kelvin one, where the required mobility is $\sim 10^4 \text{ cm}^2/\text{Vs}$ rather than $\sim 10^6$ — a growth problem, not a physics problem. The two measurements are structured predictions: each follows from the device geometry and the cited signature, and each is a pass/fail test the field can run.

The dictionary that opens this paper is, in the end, an engineering resource: strain writes gauge and metric fields into a band structure at scales — hundreds of tesla, tenths of a kelvin — that no external apparatus reaches, and it does so reproducibly enough to design with [44]. We have made that resource quantitative across named materials, drawn one concrete horizon device from it with its temperature computed from measured parameters, and stated the disorder threshold that today keeps that device from being built, as a number the field can aim at rather than a caveat it can wave away. Every step across the analogy has been graded in the sentence where it was taken, and the device lives entirely on the kinematic side of the partition its companion paper fixes [13]: it engineers the band-geometry the analogy is exact about, and it claims nothing about gravity. A strained crystal is a designer spacetime for its electrons — and the two measurements above are how we would find out whether the one we designed can be seen.

A The Painlevé–Gullstrand metric from the tilted Weyl Hamiltonian

This appendix derives the effective metric of Eq. (2), locates the horizon, and fixes the normalization of the analog surface gravity κ_{eff} used throughout the paper and in the seeded computation of Appendix C. The result is a convention, stated explicitly so that no hidden factor of order unity inflates the analog Hawking temperature.

From the Hamiltonian to a null cone. Start from the low-energy Weyl Hamiltonian with a spatially varying node shift and tilt,

$$H = v_F \boldsymbol{\sigma} \cdot (\mathbf{k} - \mathbf{b}(\mathbf{r})) + v_F W(\mathbf{r}) \cdot \mathbf{k}, \quad (6)$$

where $\mathbf{b}(\mathbf{r})$ is the strain-induced chiral (axial) gauge field [14] and $W(\mathbf{r})$ the tilt vector, type-II where $|W| > 1$ [3, 74]. The eigenvalues are

$$E_{\pm}(\mathbf{k}, \mathbf{r}) = v_F W \cdot \mathbf{k} \pm v_F |\mathbf{k} - \mathbf{b}|. \quad (7)$$

The node shift $\mathbf{b}$ enters exactly as a vector potential and can be gauged into the origin of $\mathbf{k}$ locally; it does not affect the tilt geometry, so we set $\mathbf{b} = 0$ in what follows and restore it as a gauge field where needed. Squaring

the tilt term to the other side, the dispersion is the null condition of a quadratic form,

$$(E - v_F W \cdot \mathbf{k})^2 - v_F^2 \mathbf{k}^2 = 0 \equiv g^{\mu\nu} p_\mu p_\nu = 0, \quad (8)$$

with $p_\mu = (E, \mathbf{k})$, which reads off an effective inverse metric

$$g^{00} = 1, \quad g^{0i} = -v_F W^i, \quad g^{ij} = v_F^2 (W^i W^j - \delta^{ij}). \quad (9)$$

The carriers are massless fermions propagating on the null cones of Eq. (9); this is the exact, kinematic content of the strain-to-geometry analogy [5, 63, 81].

Painlevé–Gullstrand form. The device delivers a uniaxial strain gradient, so the tilt is along the channel, $W(\mathbf{r}) = W(x) \hat{\mathbf{x}}$. Inverting Eq. (9) for this case gives, up to an overall conformal factor that affects neither the null structure nor the surface gravity, the line element

$$ds^2 = -(v_F^2 - v(x)^2) dt^2 - 2v(x) dt dx + dx^2 + dy^2 + dz^2, \quad v(x) \equiv v_F W(x), \quad (10)$$

which is the Painlevé–Gullstrand metric with “light speed” v_F and “flow velocity” $v(x)$ [5, 36]. The radial null rays follow from $ds^2 = 0$ with $dy = dz = 0$,

$$\frac{dx}{dt} = v(x) \pm v_F = v_F (W(x) \pm 1), \quad (11)$$

the same two characteristics carried by the tilted cone directly. The outgoing ray $dx/dt = v_F(W - 1)$ changes sign where $W(x) = 1$: for $W > 1$ (type-II) both rays are dragged forward and nothing escapes upstream, a trapped region; for $W < 1$ (type-I) upstream escape is allowed. The surface $W = 1$, where $v(x) = v_F$ and g^{00} changes sign, is the electronic horizon [43, 63].

The surface-gravity normalization, with its $O(1)$ factor shown. The analog surface gravity of a Painlevé–Gullstrand horizon is defined from the gradient of the lapse combination at the horizon [5],

$$\kappa_{\text{eff}} = \frac{1}{2v_F} \left| \frac{d}{dx} (v_F^2 - v^2) \right|_{W=1}. \quad (12)$$

With $v = v_F W$ one has $d(v_F^2 - v^2)/dx = -2v_F^2 W (dW/dx)$, which at the crossing $W = 1$ is $-2v_F^2 (dW/dx)$. The factor $\frac{1}{2}$, the division by v_F , and the factor of 2 from differentiating v^2 combine to unity, leaving

$$\boxed{\kappa_{\text{eff}} = v_F \left| \frac{\partial W}{\partial x} \right|_{W=1}} \quad (13)$$

with no residual numerical prefactor. The same value follows from the independent definition of the surface gravity as the peeling rate of the outgoing null ray, $\kappa_{\text{eff}} = |d(dx/dt)_{\text{out}}/dx|_{W=1} = v_F |\partial W/\partial x|$ from Eq. (11),

so the result is robust to the choice among the standard, physically equivalent $O(1)$ definitions. Equation (13) is the convention this paper adopts, and it coincides with Volovik’s tilt-gradient temperature $T_H^a = (\hbar/2\pi k_B) \partial v/\partial x$ [81] and with De Beule et al.’s $T_H^a = \hbar v_F \alpha / 2\pi k_B$ under the identification $\alpha = \partial W/\partial x$ [18]. The analog Hawking temperature is then

$$T_H^a = \frac{\hbar \kappa_{\text{eff}}}{2\pi k_B} = \frac{\hbar v_F}{2\pi k_B} \left| \frac{\partial W}{\partial x} \right|_{W=1}. \quad (14)$$

Every numerical value of κ_{eff} and T_H^a in Section 3, the kill threshold of Section 4, the signature magnitudes of Section 5, and each figure that plots them uses Eq. (13) as written; the seeded computation of Appendix C carries the same normalization, and this is the only place the $O(1)$ factor is set. The construction is exact on the kinematic side of the analog partition and silent on the dynamic side, as fixed in Section 1 [13]: Eq. (10) is a real effective metric for the carriers’ dispersion, and it is not a solution of any field equation.

B Shortlist parameter provenance, entry by entry

Every number in Table 1 is either a measured value with a primary source or a computed value flagged as computed. This appendix gives the provenance row by row, and states for each entry whether it is a measurement or a calculation, so that the kill-condition verdict of Section 4 rests on traceable inputs.

Graphene. The Fermi velocity $v_F \sim 1 \times 10^6$ m/s, the hopping $t \sim 2.7\text{--}3$ eV, and the lattice constant $a = 1.42$ Å are the standard Dirac parameters [11, 83]. The mobility entry $\mu > 2 \times 10^5$ cm²/Vs is the measured suspended-graphene value [7], with boron-nitride encapsulation delivering comparable or better mobility and micrometre-scale ballistic transport at room temperature [17, 55]. The achievable strain has two measured anchors: sample-wide uniform elastic strain to $\sim 6\%$ (peak $\sim 5.8\%$) in in-situ tensile testing [10], and an intrinsic elastic limit near $\sim 25\%$ inferred from the nanoindentation strength (42 N/m, Young’s modulus 1.0 TPa) [49]. Graphene is a two-dimensional Dirac system with no cone tilt, so its rest tilt is 0 and $\partial W/\partial u$ is undefined; it enters the shortlist as the pseudo-field champion and mobility benchmark, not as a horizon host. Its pseudo-gauge coupling $\beta_G \approx 2$ [2, 83], with the tight-binding spread $\beta_G \approx 2.6\text{--}4.0$ [66] and the distinct phonon Grüneisen parameter $\gamma(E_{2g}) \approx 1.99$ [58], anchors the Grüneisen estimate used for the tilt sensitivity below.

WTe₂. Type-II Weyl by prediction [74] and by laser-photoemission observation of Fermi arcs [85], with the caveat that the arc’s topological triviality was shown

not to confirm the tilt by itself [8]; the rest tilt $\gtrsim 1$ is therefore computed, not measured. The Fermi velocity entered is the design host value $\sim 2 \times 10^5$ m/s (Appendix C), not a separate measurement. The mobility $\sim 10^3\text{--}10^4$ cm²/Vs is the order of magnitude for the compound from its magnetotransport [1]; we do not attribute a single-digit mobility to that paper, whose two-carrier fit sits in its paywalled extended data, and the kill-condition table uses the compiled representative value of Appendix C. The achievable strain $\sim 0.23\%$ is the demonstrated cryogenic ceiling of a piezoelectric strain cell of the relevant class [32].

MoTe₂ and Mo_{0.91}W_{0.09}Te₂. Type-II Weyl by prediction [77, 84] and by photoemission [21, 34, 40]; four type-II Weyl points sit ~ 6 meV above E_F with a Weyl-pair spacing $\sim 4\%$ of the reciprocal lattice vector [77]. The strain-tunable alloy Mo_{0.91}W_{0.09}Te₂ is adopted as the concrete lead because calibrated operando strain to 0.34% (reversible below 0.17%) has been delivered on it, driving a $\sim 30\%$ increase in the density of states at E_F — but reporting no tilt parameter and locating no type-I/type-II crossing [48], so the achievable-strain entry 0.34% is measured while the rest tilt $\gtrsim 1$ is computed. Ambient-pressure superconductivity at $T_c = 0.10$ K [69] and ohmic contacts by laser phase-patterning [12] are the process anchors for the device rather than shortlist parameters.

Cd₃As₂. Included as an untilted ultrahigh-mobility anchor. The measured mobility exceeds 10^7 cm²/Vs below 4 K [51], and the Fermi velocity $\sim 1.5 \times 10^6$ m/s and lattice/torsion parameters are those used in the strain-gauge modelling of the compound [68], which also gives the modelled torsion strain window ($\sim 3\%$ unidirectional, up to $\sim 8\%$ before lattice breakdown). It is a type-I Dirac semimetal: rest tilt 0, no $\partial W/\partial u$, no horizon crossing to engineer.

NbP. Included as an untilted ultrahigh-mobility Weyl anchor of the TaAs class. The measured mobility is 5×10^6 cm²/Vs at 1.85 K, with a magnetoresistance of 850,000% at 1.85 K falling to 250% at room temperature [73]; the compound also carries the mixed axial-gravitational anomaly signature [24]. It is type-I: rest tilt < 1 , no tilt to bend.

8-Pmmn borophene (model). A tilted-Dirac cross-check rather than a fabricable host. The tight-binding model gives closed-form strain couplings for both a pseudo-magnetic vector potential and a scalar potential of the same family as the Weyl case [89]; there is no measured mobility, and the entries are marked “model.”

The tilt sensitivity $\partial W/\partial u$ (computed, all type-II rows). This is the one quantity the paper computes

rather than cites, because no experiment has measured it in any type-II material [37, 48]. Two independent routes fix its range. The density-functional strain sensitivities of the type-II literature give the first: a 0.5% compression drives the WTe₂ Weyl-point separation from $\sim 0.7\%$ to $\sim 2\%$ of the reciprocal lattice vector [74], and MoTe₂ is predicted to cross the type-II/type-I boundary under strain [77, 84]. Reading a boundary crossing $\Delta W \sim 0.1$ – 0.3 over ~ 0.5 – 2% strain places $\partial W/\partial u \sim 5$ – 30 . The second route is a Grüneisen estimate $\partial W/\partial u \sim W\beta_G$ with $\beta_G \approx 2$ – 3 (from the graphene electronic and phonon Grüneisen values above), which lands in the same band. The seeded computation of Appendix C propagates the conservative, central, and optimistic values 5, 15, 30 throughout; the factor-of-six spread they induce in T_H^a is the direct cost of this being computed rather than measured, and closing it by direct measurement is the program’s first milestone (Section 6).

C The seeded computation

The design envelope, the analog Hawking temperatures, the kill-condition thresholds, and the signature magnitudes come from one script, `sim_weyl.py`, whose source and outputs (`sim_weyl_output.txt/.json`) ship with the paper. The computation is analytic; the seed recorded in the source (20260919) exists only for provenance.

The tilt sensitivity is computed, not measured

Every strain→tilt number in this paper rests on the derivative $\partial W/\partial u$, and no experiment has measured it: the most recent operando-strain study reaches 0.34% strain on Mo_{0.91}W_{0.09}Te₂ but reports band and transport changes without a tilt parameter [48]. We therefore compute it and carry a range. The anchors are the density-functional strain sensitivities of the type-II literature — a 0.5% compression that drives the WTe₂ Weyl-point separation from $\sim 0.7\%$ to $\sim 2\%$ of the reciprocal lattice vector [74], and the DFT-predicted strain-driven type-II/type-I transitions of Sun et al. [77] and Wang et al. [84] — together with a Grüneisen estimate $\partial W/\partial u \sim W\beta_G$ with $\beta_G \approx 2$ – 3 . Reading a boundary crossing $\Delta W \sim 0.1$ – 0.3 over ~ 0.5 – 2% strain, both routes land in the band $\partial W/\partial u \sim 5$ – 30 ; the script propagates the conservative, central, and optimistic values 5, 15, 30 throughout.

Device, kill, and signatures

With a host Fermi velocity $v_F = 2 \times 10^5$ m/s and a strain of 0.3% delivered over $L = 100$ nm, the central sensitivity gives $\partial W/\partial x = 0.45 \mu\text{m}^{-1}$, $\kappa_{\text{eff}} = 9.0 \times 10^{10} \text{ s}^{-1}$, and $T_H^a = 0.11$ K (Propositions 2; Fig. 2); across the envelope (0.1–0.5% strain, 50 nm–1 μm , sensitivity 5–30)

T_H^a spans 3.6 mK–0.36 K, with κ_{eff} carrying the $O(1)$ normalization fixed in Appendix A. The kill threshold is $\mu_{\text{min}}(T_H^a) = ev_F/(k_F k_B T_H^a)$, evaluated with k_F from a representative $n \sim 10^{18} \text{ cm}^{-3}$: at the 0.11 K design point $\mu_{\text{min}} \approx 7 \times 10^5 \text{ cm}^2/\text{Vs}$, so the near-critical type-II hosts ($\mu \sim 10^3$ – 10^4) fall two orders short while Cd₃As₂ and NbP clear it but carry no tilt (Fig. 3); the threshold falls as $1/T_H^a$, reaching $\sim 10^4 \text{ cm}^2/\text{Vs}$ only near $T_H^a \sim 10$ K. The signatures (Fig. 4) are the shot-noise thermal factor at T_H^a (observable for a bath below ~ 0.1 K), the $\sim 2e^2/h$ -per-mode conductance peak of de Beule et al. [18], and a horizon dissipation of $\sim 2 \times 10^{-14} \text{ W}$, at the SQUID-on-tip readout floor [29] — which is why thermometry is the hardest of the three.

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