# Emergence of the Physical Laws of a Macroscopic Observer

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**Abstract**

We ask how apparently fundamental laws can arise from the physics of observation itself. In the Conservation-Congruent Encoding framework, an observer is a finite material device whose records must be stored and repeatedly reset. We make that bookkeeping explicit through developing an example case of a one-bit observer: a particle in a symmetric double-well potential immersed in a homogeneous thermo-acoustic medium. Incoming acoustic packets flip the bit reversibly, while a time-dependent control protocol restores the ready state. Because reset is logically irreversible and occurs while the bit remains coupled to the bath, each cycle dissipates at least $k_B T_0 \Delta\mathcal{H}_{cg}$, so under a reset rate $\nu$ the observer acts as a localized heat source with mean power $P \geq \nu k_B T_0 \Delta\mathcal{H}_{cg}$. In steady state this produces a thermal halo $\delta T(r) = P/(4\pi k r)$, which in a focusing medium induces the refractive profile $n(r) = n_0(1 + \Gamma/r)$. The resulting ray bending enlarges the capture cross-section and, in the weak-field limit, is mathematically equivalent to motion in an attractive $1/r$ potential. An external analyst restricted to the reduced event stream can therefore mistake self-induced bath distortion for an intrinsic force law. We then sketch a speculative gravitational extrapolation in which erased information is absorbed by local horizons, Newton's constant becomes Vacuum Informational Compliance, and a positive cosmological constant sets a deep-field crossover scale; with additional equilibrium assumptions, the weak-field closure can then be lifted toward the Einstein field equations.

## 1 Introduction

The historical pursuit of fundamental physical laws often relies on a "view from nowhere": an idealized observer that can register a dynamical system without materially participating in it. A physical account of information undermines that abstraction. Landauer's analysis of logical irreversibility, together with Bennett's reversible-computation perspective, makes clear that persistence, reversible transport, and erasure have distinct physical statuses [10, 3]. An observer is therefore not a mathematical ghost, but a finite material device whose records must be physically realized, stabilized, and eventually overwritten.

The Conservation-Congruent Encoding (CCE) framework sharpens this point by defining information directly in terms of coarse-grained dynamics [5]. Under a chosen coarse-graining, a macroscopic distinction qualifies as information only insofar as it remains operationally stable against ambient fluctuations. Observation is thus organized by a projection $\Pi : X \to \mathcal{Z}$ from a microscopic state space $X$ to a macroscopic state space $\mathcal{Z}$ of slow, operationally resolvable variables. The fast degrees of freedom discarded by $\Pi$ are not removed from the physics; they persist as an unresolved bath with which the retained macroscopic variables must continuously exchange conserved quantities.

Because those macroscopic distinctions must be actively protected, the projection is not passive. In the broader CCE program, one can summarize that bookkeeping by separating reversible transport from dissipative merger of logical histories [13, 14, 5]. A finite observer that continuously tracks its environment must therefore continually reuse encodings. Each reset destroys coarse-grained distinguishability and must dump the corresponding entropy into surrounding degrees of freedom. The basic question pursued here is what later regularities can be induced when that unavoidable informational exhaust feeds back into the medium through which subsequent observations are made.

This note approaches that question in the simplest fully explicit setting we could find. We construct a one-bit observer realized by a particle in a double-well potential immersed in a compressible fluid. Incoming acoustic packets flip the bit reversibly under a static Hamiltonian, while an explicit time-dependent control protocol restores the ready state and dissipates at least $k_B T_0 \Delta \mathcal{H}_{cg}$ per reset cycle, with $\Delta \mathcal{H}_{cg} = \ln 2$ only in the maximally uncertain one-bit limit. Classical heat diffusion then generates a $1/r$ thermal halo, and the temperature dependence of acoustic propagation turns that halo into a refractive lens that alters the observer's future capture cross-section. The first half of the paper therefore isolates, in a controlled toy model, a concrete mechanism by which observer maintenance can masquerade as an external attractive law.

That toy model is valuable chiefly because of what it enables in the latter sections. Once the thermo-acoustic mechanism has been made explicit, the discussion shifts from the example itself to its intended interpretation: whether the same CCE bookkeeping can be read more generally when the substrate absorbing erased information is no longer an ordinary fluid but the local vacuum sector relevant to gravity. In that re-reading, the weak-field $1/r$ closure becomes a prototype for observer-induced attraction, Newton's constant is recast as Vacuum Informational Compliance, local horizon balance supplies the route toward a Jacobson-style tensorial closure, and a positive cosmological constant sets the deep-field crossover scale. Those later sections are therefore not a detachable speculation appended to a toy model; they are the broader target of the note. The Single-Bit Capture Model is introduced as the cleanest concrete bridge from irreversible reset to the possibility that apparently fundamental law-like structure emerges from the material architecture of observation itself.

## 2 The Single-Bit Capture Model

To sharpen the argument, we reduce the observer to the smallest reusable register that can still support an event: a single bit. The model has four ingredients: a microscopic double-well memory, a reversible thermo-acoustic capture interaction, an explicit control protocol that resets the memory while dissipating heat, and a steady thermal closure that feeds that dissipation back into future signal propagation.

### 2.1 Microscopic Bit, Fluid Bath, and Coarse-Graining

We model the surrounding bath as a homogeneous compressible fluid occupying $\mathbb{R}^3$. Its fast acoustic degrees of freedom are represented by a scalar field $\phi(x,t)$ and conjugate momentum $\pi_\phi(x,t)$, while the coarse temperature field $T(x,t)$ is the corresponding local temperature after coarse-graining over those microscopic bath modes. A generic microscopic state is therefore

$$(q, p, \phi, \pi_\phi) \in X, \tag{1}$$

where $(q,p)$ are the observer's encoding coordinates. The bit itself is realized by a particle of mass $m$ in a controllable double-well potential

$$V(q,t) = \frac{\lambda}{4} q^4 - \frac{\mu(t)}{2} q^2 - f(t) q, \qquad \lambda > 0. \tag{2}$$

During idle storage and signal capture the control parameters are held fixed at

$$\mu(t) = \mu_0 > 0, \qquad f(t) = 0, \tag{3}$$

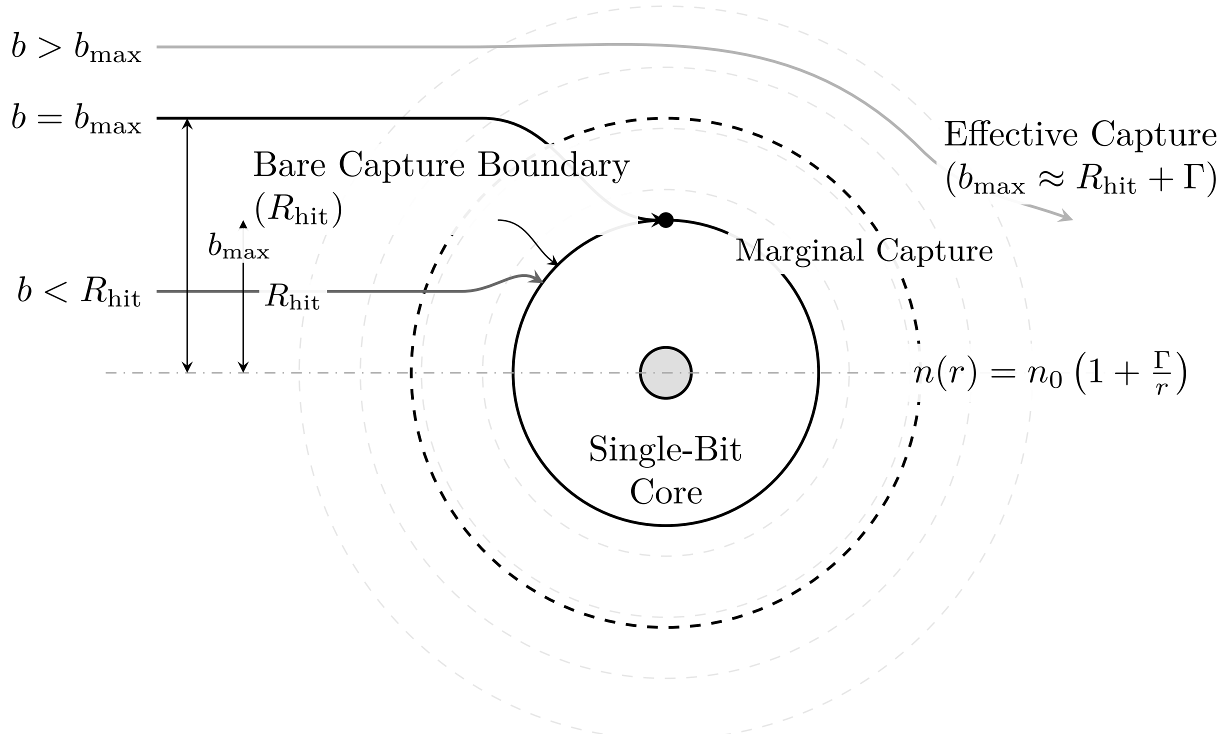


Figure 1: The effective capture cross-section of a single-bit observer. Here $R_{\text{hit}}$ denotes the bare capture radius in the reduced ray model; the effective source radius of the time-averaged thermal exhaust is treated separately in the text. A time-dependent reset protocol erases the bit and dissipates heat into the surrounding fluid. In a focusing medium with positive thermo-acoustic coefficient $(\partial n/\partial T)_{T_0} > 0$, the resulting steady $1/r$ thermal halo establishes a refractive lens. Incident signals with impact parameters $R_{\text{hit}} < b \leq b_{\max}$ that would bypass the observer in flat space are deflected inward and registered. An external analyst therefore measures an augmented cross-section ($\sigma_{\text{eff}} > \sigma_0$), inferring the presence of an emergent attractive source under the stated constitutive assumptions.

so the memory sees the static symmetric potential

$$V_0(q) = \frac{\lambda}{4}q^4 - \frac{\mu_0}{2}q^2. \tag{4}$$

This potential has two stable minima,

$$q_\pm = \pm q_0, \qquad q_0 = \sqrt{\mu_0/\lambda}, \tag{5}$$

separated by the barrier height

$$\Delta V = V_0(0) - V_0(q_\pm) = \frac{\mu_0^2}{4\lambda}. \tag{6}$$

We take the left well to be the ready state and the right well to be the registered-hit state,

$$0 \leftrightarrow q \approx -q_0, \qquad 1 \leftrightarrow q \approx +q_0. \tag{7}$$

To make the carrier medium explicit, we choose the linear acoustic bath Hamiltonian

$$H_{\text{bath}}[\phi, \pi_\phi] = \frac{1}{2} \int_{\mathbb{R}^3} \left( \frac{\pi_\phi(x)^2}{\rho_0} + \rho_0 c_0^2 |\nabla \phi(x)|^2 \right) d^3x, \tag{8}$$

where $\rho_0$ is the background mass density and $c_0$ is the ambient sound speed. The bit couples locally to the fluid through the total Hamiltonian

$$H_{\text{tot}}(t) = \frac{p^2}{2m} + V(q,t) + H_{\text{bath}}[\phi, \pi_\phi] - \gamma q\, \phi(x_{\text{obs}}), \tag{9}$$

with coupling strength $\gamma$ and observer location $x_{\text{obs}}$. The observer does not retain the full microstate. Its operational state space is the coarse-grained bit

$$\Pi : X \to \mathcal{Z} = \{0, 1\}, \qquad \Pi(q, p, \phi, \pi_\phi) = \begin{cases} 0, & q < 0, \\ 1, & q > 0. \end{cases} \tag{10}$$

The dividing surface $q = 0$ is unstable and may be assigned arbitrarily at measure zero. This is the core simplification of the model: the observer stores only one reusable bit, while all bath detail beyond that sign is discarded.

## 2.2 Reversible Capture

During capture the control parameters remain frozen at $\mu(t) = \mu_0$ and $f(t) = 0$, so $V(q,t) = V_0(q)$ and the full dynamics are generated by a time-independent Hamiltonian. For the particle degree of freedom, the canonical equations follow from ordinary derivatives of $H_{\text{tot}}$:

$$\dot{q} = \frac{\partial H_{\text{tot}}}{\partial p} = \frac{p}{m}, \qquad \dot{p} = -\frac{\partial H_{\text{tot}}}{\partial q} = -V_0'(q) + \gamma\, \phi(x_{\text{obs}}, t). \tag{11}$$

For the continuous acoustic bath, Hamilton's equations are written in terms of functional derivatives,

$$\partial_t \phi(x,t) = \frac{\delta H_{\text{tot}}}{\delta \pi_\phi(x,t)}, \qquad \partial_t \pi_\phi(x,t) = -\frac{\delta H_{\text{tot}}}{\delta \phi(x,t)}. \tag{12}$$

To expose the source term, it is convenient to rewrite the interaction energy as

$$H_{\text{int}} = -\gamma q(t) \int_{\mathbb{R}^3} \delta^{(3)}(x - x_{\text{obs}})\, \phi(x,t)\, d^3x. \tag{13}$$

Using this form, and integrating the bath term by parts with vanishing boundary contribution at infinity, gives

$$\frac{\delta H_{\text{tot}}}{\delta \pi_\phi(x,t)} = \frac{\pi_\phi(x,t)}{\rho_0}, \qquad \frac{\delta H_{\text{tot}}}{\delta \phi(x,t)} = -\rho_0 c_0^2 \nabla^2 \phi(x,t) - \gamma q(t)\, \delta^{(3)}(x - x_{\text{obs}}). \tag{14}$$

Hence the bath satisfies the explicit canonical field equations

$$\partial_t \phi(x,t) = \frac{\pi_\phi(x,t)}{\rho_0}, \qquad \partial_t \pi_\phi(x,t) = \rho_0 c_0^2 \nabla^2 \phi(x,t) + \gamma q(t)\, \delta^{(3)}(x - x_{\text{obs}}). \tag{15}$$

Taking a time derivative of the first field equation and substituting the second yields the classical sourced wave equation

$$\partial_t^2 \phi - c_0^2 \nabla^2 \phi = \frac{\gamma}{\rho_0} q(t)\, \delta^{(3)}(x - x_{\text{obs}}). \tag{16}$$

Because this is the full field equation, the point value $\phi(x_{\text{obs}}, t)$ contains both the incoming packet and the observer's own retarded self-field. To isolate the externally delivered write work, we therefore decompose the local field during the capture window as

$$\phi(x,t) = \phi_{\text{in}}(x,t) + \phi_{\text{self}}(x,t), \tag{17}$$

where $\phi_{\text{in}}$ denotes the incident wave packet that would be present in the absence of observer backreaction over that short interaction interval. Any static self-contribution may then be absorbed into a renormalization of the trapping landscape or, in a finite-size regularization, into the local coupling kernel. The mechanical work transferred to the bit by the incoming packet is therefore

$$W_{\text{in}} = \int_{t_i}^{t_f} \gamma\, \phi_{\text{in}}(x_{\text{obs}}, t)\, \dot{q}(t)\, dt. \tag{18}$$

If $W_{\text{in}}$ exceeds the barrier $\Delta V$ and the post-encounter relaxation leaves the particle trapped in the right metastable basin, the trajectory crosses $q = 0$ and the coarse state durably flips from 0 to 1. Because the control parameters are fixed and the bath is included explicitly, this capture step is reversible Hamiltonian transport on the full phase space: Liouville volume is preserved and no logical information is erased, even though coarse-grained relaxation can stabilize the write within a single logical basin.

Because the reduced interaction in $H_{\text{tot}}$ remains pointlike, the model above does not by itself define a finite geometric cross-section. We therefore introduce a bare hit radius $R_{\text{hit}}$: in a regularized finite-size interaction model it is the largest impact parameter for which the incident work still satisfies $W_{\text{in}} \geq \Delta V$. In the later eikonal treatment, the marginal captured ray is taken to graze the corresponding effective capture surface at $r = R_{\text{hit}}$. The effective support radius of the time-averaged thermal source will be denoted separately by $R_{\text{src}}$. In a more microscopic treatment one would derive both scales from the packet profile, coupling kernel, and threshold dynamics; here $R_{\text{hit}}$ parametrizes the regularized bare capture threshold while $R_{\text{src}}$ parametrizes the heat-emitting core.

## 2.3 Controlled Reset and Dissipation

A one-bit observer cannot continue registering events unless the bit is reused. The reset map is therefore the logically irreversible operation

$$0 \mapsto 0, \qquad 1 \mapsto 0. \tag{19}$$

We implement reset by an explicit control protocol on the potential $V(q,t)$. A minimal erase cycle proceeds in three stages: first lower the barrier by decreasing $\mu(t)$, then tilt the potential toward the ready well by taking $f(t) < 0$, and finally restore the original double-well by raising $\mu(t)$ back to $\mu_0$ and removing the tilt after the particle has relaxed into the left basin. Because $V$ now depends explicitly on time, the controller supplies work

$$W_{\text{ctrl}} = \int_{t_i}^{t_f} \frac{\partial V}{\partial t}(q(t), t)\, dt. \tag{20}$$

The bit remains coupled to the fluid throughout this protocol. After coarse-graining over bath modes on times longer than the acoustic correlation time, the controlled degree of freedom obeys the effective Langevin equation

$$m\ddot{q} + \eta\dot{q} + \partial_q V(q,t) = \zeta(t), \qquad \langle \zeta(t)\zeta(t') \rangle = 2\eta k_B T_0\, \delta(t - t'), \tag{21}$$

where $\eta$ is the effective damping coefficient and $T_0$ is the ambient bath temperature. The term $\eta\dot{q}$ is the explicit mechanical channel by which control work is dumped into the fluid. Let $(p_0, p_1)$ denote the logical occupancies sampled immediately before a blind reset clock fires. To make those occupancies explicit for the one-bit observer, suppose an externally driven unconditional reset clock of period $\tau = 1/\nu$ acts while successful captures arrive as a Poisson process with effective rate $\lambda_{\text{hit}}$. Because the register can store at most one unresolved hit between resets, only the first capture in a cycle changes the logical state, so

$$p_0 = e^{-\lambda_{\text{hit}}\tau}, \qquad p_1 = 1 - e^{-\lambda_{\text{hit}}\tau}. \tag{22}$$

The erased coarse-grained entropy per cycle is then

$$\Delta\mathcal{H}_{cg} = -p_0 \ln p_0 - p_1 \ln p_1 = h_2(p_1) \le \ln 2, \qquad h_2(u) := -u \ln u - (1-u)\ln(1-u), \tag{23}$$

with the upper bound attained only for a maximally uncertain one-bit state. Since the protocol maps two metastable logical alternatives onto one, Landauer's bound applies: for any erase executed against a bath at temperature $T_0$,

$$\langle Q_{\text{bath}} \rangle \ge k_B T_0 \Delta\mathcal{H}_{cg}, \tag{24}$$

with equality only in the quasistatic limit [10, 4, 9]. At reset rate $\nu$, the observer therefore acts as a localized heat source of mean power

$$P := \nu \langle Q_{\text{bath}} \rangle \ge \nu k_B T_0 \Delta\mathcal{H}_{cg}. \tag{25}$$

In the special case $p_0 = p_1 = 1/2$, this reduces to $P \ge \nu k_B T_0 \ln 2$. The general form above will be needed below once the reset clock is tied to the hit process and $\Delta\mathcal{H}_{cg}$ is no longer an independent input.

At the level of coarse observables, this explicit machine also clarifies how the toy model fits into the broader CCE language. After eliminating the fast acoustic bath modes, one may summarize the effective macroscopic dynamics schematically by

$$\dot{A} = \{A, H_{\text{eff}}(t)\} + (A, S_{\text{eff}}). \tag{26}$$

During the capture phase, with the controls frozen, the reversible acoustic kick is generated by the Hamiltonian Poisson sector $\{A, H_{\text{eff}}\}$, representing reversible, energy-conserving transport. During reset, the externally driven potential together with the dissipative term $\eta\dot{q}$ contracts the two logical basins onto the ready basin while exporting entropy to the fluid; in the CCE formalism, that coarse-grained merger is summarized by the symmetric metric sector $(A, S_{\text{eff}})$, representing dissipative, entropy-producing merger. The metriplectic split is therefore not the microscopic mechanism in this toy model; it is the effective macroscopic summary of the explicit thermo-acoustic machine that implements reversible capture and irreversible erasure.

## 2.4 Backreaction and Emergent Capture Law

We now close the toy model by following the fate of that dissipated heat in the same medium that carries the incoming signals. Let $R_{\text{src}}$ denote the effective support radius of the time-averaged heat source created by repeated resets. Assuming a conduction-dominated steady state over an observation time $t \gg \max(\nu^{-1}, R_{\text{src}}^2/D_T)$, where $D_T$ is the fluid's thermal diffusivity, the repeated erasures may be replaced by a continuous mean heat source of total power $P$ supported inside $r \leq R_{\mathrm{src}}$. Outside that compact source region, the steady temperature perturbation

$$\delta T(r) = T(r) - T_0 \tag{27}$$

obeys the sourceless heat equation

$$\nabla^2 \delta T = 0, \qquad r > R_{\mathrm{src}}, \tag{28}$$

with asymptotic condition $\delta T(r) \to 0$ as $r \to \infty$ and total outward heat flux fixed by Fourier's law,

$$-4\pi r^2 k \, \frac{d\delta T}{dr} = P, \qquad r > R_{\mathrm{src}}, \tag{29}$$

where $k$ is the thermal conductivity of the fluid. The unique decaying spherically symmetric solution is therefore

$$\delta T(r) = \frac{P}{4\pi k r}, \qquad r > R_{\mathrm{src}}. \tag{30}$$

The $1/r$ profile is thus not assumed; it is the exact exterior solution of steady heat diffusion in three dimensions outside the heat-emitting core.

The incoming signals are acoustic waves, so their local propagation speed is the temperature-dependent sound speed $c(T)$. Writing the refractive index as

$$n(r) = \frac{c_0}{c(T(r))}, \qquad c_0 := c(T_0), \qquad n_0 := n(T_0) = 1, \tag{31}$$

and linearizing around the ambient state gives

$$n(r) = n_0 \left[1 + \beta_T \, \delta T(r)\right] + \mathcal{O}\Big(\delta T(r)^2\Big), \qquad \beta_T := \frac{1}{n_0} \left(\frac{\partial n}{\partial T}\right)_{T_0}. \tag{32}$$

Substituting the exact thermal profile yields

$$n(r) = n_0 \left(1 + \frac{\Gamma}{r}\right), \qquad \Gamma = \beta_T \frac{P}{4\pi k}. \tag{33}$$

Enhanced capture requires $\Gamma > 0$, equivalently $\beta_T > 0$. If $\beta_T < 0$ instead, the same mechanism produces defocusing rather than attraction, so the sign of the thermo-acoustic coefficient is physically decisive. For an ideal gas, $c(T) \propto T^{1/2}$ and therefore $\beta_T = -1/(2T_0) < 0$. A focusing realization of the toy model therefore requires a different thermo-acoustic medium satisfying $\partial c/\partial T < 0$ over the relevant operating regime; identifying a concrete material and frequency window is left for future work. Until such a medium is specified, the attractive branch should be read as conditional rather than as a material-specific prediction.

To pass from the wave description to a capture cross-section, we now restrict attention to a geometric-acoustics regime in which the incident packets are narrow-band and their wavelength is short compared with both $R_{\mathrm{hit}}$ and the scale over which $n(r)$ varies. In that eikonal limit, propagation may be treated in terms of rays obeying Fermat's principle. We assume the effective capture surface lies outside or near the edge of the heated core, $R_{\mathrm{src}} \lesssim R_{\mathrm{hit}}$, so that the exterior refractive profile can be evaluated at the marginal trajectory.

In the weak-lensing regime $|\Gamma| \ll b$, a ray arriving from infinity with impact parameter $b$ experiences the signed deflection

$$\Delta\theta(b) \approx -\frac{1}{n_0} \int_{-\infty}^{\infty} \frac{\partial n}{\partial b} \, dx = \Gamma \int_{-\infty}^{\infty} \frac{b \, dx}{(x^2 + b^2)^{3/2}} = \frac{2\Gamma}{b}. \tag{34}$$

Without thermal backreaction, capture requires $b \le R_{\rm hit}$, so the bare cross-section is

$$\sigma_0 = \pi R_{\rm hit}^2. \tag{35}$$

For the spherically symmetric refractive profile above, Fermat's principle gives the conserved optical angular momentum

$$L = n(r)\, r \sin\psi, \tag{36}$$

where $\psi$ is the local propagation angle relative to the radial direction. A ray arriving from infinity has $L = n_0 b$, while the marginal captured ray just grazes the effective capture surface at $r = R_{\rm hit}$ with $\sin\psi = 1$. In the focusing regime $\Gamma > 0$, conservation of $L$ therefore gives the exact threshold within the reduced eikonal model

$$b_{\max} = \frac{n(R_{\rm hit})}{n_0} R_{\rm hit} = R_{\rm hit}\left(1 + \frac{\Gamma}{R_{\rm hit}}\right) = R_{\rm hit} + \Gamma, \tag{37}$$

and hence

$$\sigma_{\rm eff} \approx \pi\,(R_{\rm hit} + \Gamma)^2 = \pi R_{\rm hit}^2 + 2\pi R_{\rm hit}\Gamma + \mathcal{O}(\Gamma^2). \tag{38}$$

Combining this with the Landauer bound on $P$ makes the dependence on irreversible throughput explicit:

$$\Gamma \ge \beta_T \frac{\nu k_B T_0 \Delta\mathcal{H}_{cg}}{4\pi k}, \qquad (\beta_T > 0). \tag{39}$$

If the ambient incident flux is $J$, then the geometric-acoustics capture model gives the effective hit rate

$$\lambda_{\rm hit} = J\sigma_{\rm eff}. \tag{40}$$

With the unconditional reset clock of period $\tau = 1/\nu$ introduced above, the pre-reset occupancies become

$$p_0 = e^{-J\sigma_{\rm eff}/\nu}, \qquad p_1 = 1 - e^{-J\sigma_{\rm eff}/\nu}, \tag{41}$$

so

$$\Delta\mathcal{H}_{cg} = h_2\Big(1 - e^{-J\sigma_{\rm eff}/\nu}\Big). \tag{42}$$

A useful analytically closed operating branch is obtained by an externally tuned matched-clock policy: an auxiliary controller, not counted as part of the one-bit memory itself, sets the reset period equal to an estimate of the mean inter-hit time, $\tau \approx \lambda_{\rm hit}^{-1}$, equivalently

$$\nu \approx J\sigma_{\rm eff}. \tag{43}$$

This branch should therefore be read as a controlled operating regime rather than as an endogenous timing capability of the bare single-bit observer. On that branch,

$$p_0 = e^{-1}, \qquad p_1 = 1 - e^{-1}, \qquad \Delta\mathcal{H}_{cg} = h_*, \tag{44}$$

with

$$h_* := h_2(1 - e^{-1}) \approx 0.658. \tag{45}$$

The leading correction to the bare cross-section is then bounded below by

$$\sigma_{\rm eff} - \sigma_0 \gtrsim \frac{\beta_T R_{\rm hit} J \sigma_{\rm eff} k_B T_0 h_*}{2k} \tag{46}$$

so

$$\sigma_{\rm eff} \gtrsim \sigma_0 + \Lambda_* J\sigma_{\rm eff}, \qquad \Lambda_* := \frac{\beta_T R_{\rm hit} k_B T_0 h_*}{2k}, \tag{47}$$

and, on the near-minimal dissipation branch where the Landauer bound is approximately saturated,

$$\sigma_{\text{eff}} \approx \frac{\sigma_0}{1 - \Lambda_* J}. \tag{48}$$

This closure is only valid while the approximations used to derive it remain satisfied, in particular

$$|\Gamma| \ll R_{\text{hit}}, \qquad |\beta_T\, \delta T(R_{\text{hit}})| \ll 1. \tag{49}$$

The formal pole at $1 - \Lambda_* J = 0$ therefore signals the breakdown of the weak-lensing and linear-response approximations rather than a controlled prediction of a literal runaway singularity. A different reset policy would lead to the same structure with $h_*$ replaced by the corresponding value of $\Delta\mathcal{H}_{cg}$. An external analyst who sees only the enhanced hit rate can still misread that self-induced bath distortion as an intrinsic attractive property of the observer.

Beyond the static capture cross-section, the spatial geometry of the signal trajectories can mimic a fundamental force at leading order. By the optical-mechanical analogy, Fermat's principle for an acoustic ray in an index $n(r)$ is mathematically equivalent to Maupertuis' principle for a classical particle with momentum $p(r) \propto n(r)$. Matching the acoustic index $n(r) = n_0(1 + \Gamma/r)$ to the non-relativistic momentum profile $p(r) \approx p_0\,(1 - U(r)/(m_{\text{test}} v_0^2))$ shows that, to leading order in $\Gamma/r$, the ray paths coincide with those of a test mass $m_{\text{test}}$ with asymptotic velocity $v_0$ in a Newtonian gravitational potential $U(r) = -m_{\text{test}} v_0^2 \Gamma/r$. An external analyst tracking the spatial ray paths will therefore infer, at the same order, an effective inward radial acceleration of $a_r \approx -v_0^2\Gamma/r^2$. The same thermodynamic exhaust that enlarges $\sigma_{\text{eff}}$ thus also projects the leading-order geometry of an inverse-square attractive field.

## 3 Toward a Gravitational Reading of the CCE Mechanism

The thermo-acoustic model isolates a general chain rather than a uniquely acoustic phenomenon. A finite observer coarse-grains a microscopic bath into a reduced operational state, preserves that state only by continual irreversible resetting, and therefore exports a conserved load into its surrounding substrate. If that substrate is three-dimensional and responds linearly in the weak-field regime, then the resulting exterior geometry reproduces the same $1/r$ potential structure as Newtonian gravity. The purpose of this section is to state that conditional emergence claim explicitly and, following the broader CCE program [5], to indicate what extra assumptions are needed to elevate the scalar weak-field closure into a geometric equation of state. The weak-field transport argument is the part directly supported by the present toy model; the later horizon-based subsections are an extrapolation beyond the thermo-acoustic channel.

### 3.1 Formal Setup

Consider a localized macroscopic observer centered at the origin. Let

$$\Pi : X \to \mathcal{Z} \tag{50}$$

be its operative coarse-graining from a microscopic state space $X$ to a finite macroscopic state space $\mathcal{Z}$. Let $\Delta\mathcal{H}_{cg}$ be the coarse-grained entropy erased per reset, let $\mathcal{F}_c$ be the reset throughput, and let $\alpha_c$ be the channel conversion factor that converts erased coarse-grained entropy into the exported conserved load of the substrate. The generalized CCE bookkeeping bound is then

$$\mathcal{J}_{\text{obs}} \ge \mathcal{J}_{c,\min} := (\mathcal{F}_c \alpha_c)\, \Delta\mathcal{H}_{cg}, \tag{51}$$

where $\mathcal{J}_{\rm obs}$ is the actual outward load current carried by the substrate. For the purposes of this section, (51) represents a fundamental CCE requirement, with the variables generalized here for mathematical simplicity. It mandates that each reset must export at least $\alpha_c \Delta\mathcal{H}_{cg}$ of a conserved substrate load at a throughput $\mathcal{F}_c$. In the thermal one-bit channel studied above, $\mathcal{J}_{\rm obs}$ is simply the emitted power $P$, the conversion factor is $\alpha_c = k_B T_0$, and (51) reduces exactly to $P \geq \nu k_B T_0 \Delta\mathcal{H}_{cg}$. In the present extrapolation, the same structure is abstracted to highlight the universal mechanism: a macroscopic observer remains classically definite only by continuously exporting a nonzero irreversible bookkeeping load into its environment, a mandate that holds regardless of which specific conserved quantity is multiplexed to carry that exhaust. The gravitational reading below then adds only three further ingredients: a linear transport law for that load, a weak-field constitutive relation between the transported load and an effective potential, and a local Clausius/horizon-balance assumption with area-per-nat parameter $A_{\rm nat}$. No other machinery from the broader CCE program is used explicitly in the derivation sketched here.

Assume next that this load is transported by a homogeneous three-dimensional substrate with effective transport coefficient $\kappa > 0$. For a stationary pointlike observer, let $\Psi(x)$ denote the corresponding substrate-load field. The static closure is

$$-\kappa \nabla^2 \Psi(x) = \mathcal{J}_{\rm obs}\, \delta^{(3)}(x). \tag{52}$$

Outside the source, (52) reduces to Laplace's equation. By spherical symmetry,

$$\Psi(r) = \frac{\mathcal{J}_{\rm obs}}{4\pi\kappa r}, \qquad r > 0. \tag{53}$$

The $1/r$ profile is therefore not imposed by hand. It is the unique exterior solution of steady transport from a localized irreversible source in three dimensions.

### 3.2 Observer-Induced Weak-Field Gravity

To connect the substrate load to geometry, assume a linear weak-field constitutive relation

$$\Phi(r) = -\xi\Psi(r), \qquad \xi > 0, \tag{54}$$

where $\Phi$ is the effective Newtonian potential reconstructed by an external analyst and $\xi$ is the substrate susceptibility. Substituting (53) into (54) gives

$$\Phi(r) = -\frac{\xi\mathcal{J}_{\rm obs}}{4\pi\kappa r} = -\frac{\mu_{\rm obs}}{r}, \qquad \mu_{\rm obs} := \frac{\xi\mathcal{J}_{\rm obs}}{4\pi\kappa}. \tag{55}$$

Hence the effective gravitational parameter obeys the lower bound

$$\mu_{\rm obs} \geq \frac{\xi}{4\pi\kappa}(\mathcal{F}_c\alpha_c)\,\Delta\mathcal{H}_{cg}. \tag{56}$$

If one now defines the externally inferred source mass by

$$M_{\rm eff} := \frac{\mu_{\rm obs}}{G}, \tag{57}$$

then

$$M_{\rm eff} = \frac{\xi\mathcal{J}_{\rm obs}}{4\pi\kappa G} \geq \frac{\xi}{4\pi\kappa G}(\mathcal{F}_c\alpha_c)\,\Delta\mathcal{H}_{cg}. \tag{58}$$

On the near-minimal branch, the mass inferred by an external analyst is therefore proportional to the observer's irreversible throughput. In the specific thermal one-bit realization,

$$M_{\text{eff}} \approx \frac{\xi}{4\pi\kappa G}\,\nu k_B T_0 \Delta\mathcal{H}_{cg}. \tag{59}$$

Within this extrapolation, what is operationally identified as rest mass is the coarse-grained signature of the rate at which the system must erase microscopic correlations to remain macroscopically definite.

**Proposition.** Under the transport law (52) and the linear response (54), the exterior coarse-grained field generated by a localized observer is indistinguishable, in the weak-field regime, from the Newtonian field of a gravitating source with parameter $\mu_{\text{obs}}$.

*Demonstration.* From (55), the acceleration of a nonrelativistic test body is

$$\mathbf{a}(r) = -\nabla\Phi(r) = -\frac{\mu_{\text{obs}}}{r^2}\,\hat{\mathbf{r}} = -\frac{GM_{\text{eff}}}{r^2}\,\hat{\mathbf{r}}. \tag{60}$$

This is exactly the Newtonian inverse-square law.

### 3.3 Optical Form and Observer Effect

The same conclusion appears in the propagation of light. In isotropic coordinates, the weak-field line element can be written as [12]

$$ds^2 = -\left(1 + \frac{2\Phi}{c^2}\right)c^2dt^2 + \left(1 - \frac{2\Phi}{c^2}\right)\left(dr^2 + r^2 d\Omega^2\right). \tag{61}$$

For null propagation, Fermat's principle yields the effective refractive index

$$n_{\text{eff}}(r) \approx 1 - \frac{2\Phi(r)}{c^2} = 1 + \frac{2\mu_{\text{obs}}}{c^2 r}. \tag{62}$$

After identifying $\mu_{\text{obs}} = GM_{\text{eff}}$, this becomes

$$n_{\text{eff}}(r) \approx 1 + \frac{2GM_{\text{eff}}}{c^2 r}, \tag{63}$$

which is precisely the weak-field optical form of Schwarzschild propagation. Thus both matter trajectories and light rays are focused by the same observer-generated $1/r$ field. An analyst who has access only to those trajectories, and not to the microscopic CCE bookkeeping that sustains the observer's classical state, is naturally led to reconstruct an attractive source. In that restricted description, gravity appears as an observer-induced effective field generated by continual irreversible coarse-graining.

This completes the weak-field part of the argument. The remaining questions are whether the conversion factor $G$ can itself be read off from the substrate and whether the scalar transport closure can be lifted to tensorial spacetime dynamics.

### 3.4 Microscopic Origin of $G$ as Vacuum Informational Compliance

The weak-field closure above still leaves $G$ as an externally inserted conversion constant. To remove that opacity, the broader CCE picture suggests one further constitutive step: the vacuum itself is treated as the substrate that absorbs irreversible observation. At the exact microscopic level, a closed Hamiltonian description preserves phase-space volume, so nothing there by itself singles out an area law. The Bekenstein-Hawking scaling must therefore arise only after coarse-graining collapses a volumetric set of microhistories into macroscopically stable equivalence classes [2, 7, 5]. In that reading, a local causal horizon is the geometric interface on which erased information is deposited.

Let $A_{\rm nat}$ denote, in the nat convention used throughout this note, the microscopic horizon area required to absorb one nat of irreversibly erased information. Then

$$\delta A = A_{\rm nat}\,\delta I_{\rm irr}. \tag{64}$$

For a local Rindler observer, the gravitational channel factor is set by the Unruh temperature [15],

$$T_U = \frac{\hbar\kappa_{\rm surf}}{2\pi k_B c}, \tag{65}$$

so the minimal per-event Landauer exhaust is

$$\delta Q = k_B T_U\,\delta I_{\rm irr} = \frac{\hbar\kappa_{\rm surf}}{2\pi c}\,\delta I_{\rm irr}. \tag{66}$$

The same local horizon obeys the first law of horizon mechanics [1, 8],

$$\delta E = \frac{\kappa_{\rm surf}c^2}{8\pi G}\,\delta A. \tag{67}$$

If the horizon is precisely the channel that absorbs the irreversible exhaust, the minimal deformation occurs when $\delta Q = \delta E$. Substituting (64) and (66) into (67) yields

$$A_{\rm nat} = \frac{4G\hbar}{c^3} = 4\ell_P^2, \tag{68}$$

and therefore

$$G = \frac{c^3 A_{\rm nat}}{4\hbar}. \tag{69}$$

In this extrapolation, Newton's constant is not a primitive force parameter. It is the macroscopic compliance parameter of the vacuum's informational capacity: the conversion factor between irreversible coarse-grained information and the area that must open up on a causal boundary to store it. Equivalently, the Bekenstein-Hawking entropy is recovered as

$$I_{\rm max} = \frac{A}{A_{\rm nat}} = \frac{A}{4\ell_P^2}, \qquad S_{\rm BH} = k_B I_{\rm max}. \tag{70}$$

### 3.5 Local Horizon Balance and Tensorial Closure

The remaining step is to replace the scalar load $\mathcal{J}_{\rm obs}$ by a relativistic exhaust description. In the metriplectic language, the local irreversible production rate is generated by the positive semi-definite metric sector, so after coarse-graining the observer's continuous reset burden appears as a relativistic load tensor $L^{(\mathrm{exh})}_{\mu\nu}$ [5]. We normalize it so that the ordinary stress-energy tensor is

$$T^{(\mathrm{exh})}_{\mu\nu} := c^2 L^{(\mathrm{exh})}_{\mu\nu}. \tag{71}$$

In the thermo-acoustic toy model this object reduces effectively to a scalar power source; in a relativistic substrate it must encode the full four-momentum flux of the exported load. With the normalization in Eq. (71), the local Clausius flux is written in terms of $L^{(\mathrm{exh})}_{\mu\nu}$, while $T^{(\mathrm{exh})}_{\mu\nu}$ is the corresponding ordinary stress-energy tensor.

Consider any spacetime point and choose a local Rindler horizon $\mathcal{H}$ through it, generated by null vectors $k^\mu$. Near that patch the approximate boost Killing field is $\chi^\mu \approx -\kappa_{\mathrm{surf}}\lambda k^\mu$, where $\lambda$ is an affine parameter along the horizon generators and $d\Sigma^\nu = k^\nu d\lambda\, dA$. The heat flux of CCE exhaust through the horizon is then [8]

$$\delta Q = \int_{\mathcal{H}} L^{(\mathrm{exh})}_{\mu\nu} \chi^\mu d\Sigma^\nu = \kappa_{\mathrm{surf}} \int dA \int_{-\epsilon}^{0} L^{(\mathrm{exh})}_{\mu\nu} k^\mu k^\nu (-\lambda)\, d\lambda. \tag{72}$$

The corresponding area response is governed by the Raychaudhuri equation for the null congruence,

$$\frac{d\theta}{d\lambda} = -\frac{1}{2}\theta^2 - \sigma_{\mu\nu}\sigma^{\mu\nu} - R_{\mu\nu}k^\mu k^\nu. \tag{73}$$

If the local vacuum is taken as instantaneously unperturbed at $\lambda = 0$, so that $\theta(0) = 0$ and $\sigma_{\mu\nu}(0) = 0$, then to first order in the affine distance

$$\theta(\lambda) = -\lambda R_{\mu\nu}k^\mu k^\nu, \tag{74}$$

and hence the area change is

$$\delta A = \int dA \int_{-\epsilon}^{0} R_{\mu\nu}k^\mu k^\nu (-\lambda)\, d\lambda. \tag{75}$$

Imposing the same CCE equation of state as above, namely that the local geometry deforms just enough to absorb the Landauer exhaust,

$$\delta Q = \frac{\kappa_{\mathrm{surf}} c^2}{8\pi G}\, \delta A, \tag{76}$$

gives

$$L^{(\mathrm{exh})}_{\mu\nu} k^\mu k^\nu = \frac{c^2}{8\pi G} R_{\mu\nu}k^\mu k^\nu. \tag{77}$$

and therefore, by (71),

$$T^{(\mathrm{exh})}_{\mu\nu} k^\mu k^\nu = \frac{c^4}{8\pi G} R_{\mu\nu}k^\mu k^\nu. \tag{78}$$

Because (78) must hold for every local observer and every null direction, one obtains

$$R_{\mu\nu} + f g_{\mu\nu} = \frac{8\pi G}{c^4} T^{(\mathrm{exh})}_{\mu\nu} \tag{79}$$

for some scalar field $f$. Requiring local conservation

$$\nabla^\mu T^{(\mathrm{exh})}_{\mu\nu} = 0 \tag{80}$$

and using the contracted Bianchi identity fixes

$$f = -\frac{1}{2}R + \Lambda, \tag{81}$$

so the geometric closure becomes

$$R_{\mu\nu} - \frac{1}{2}Rg_{\mu\nu} + \Lambda g_{\mu\nu} = \frac{8\pi G}{c^4}T^{(\mathrm{exh})}_{\mu\nu}. \tag{82}$$

Within this stronger extrapolation, the Einstein field equations are read as the macroscopic equation of state that continuously balances observer-generated CCE exhaust against local Vacuum Informational Compliance. Their nonlinearity is then not inserted separately; it is inherited from the Raychaudhuri dynamics and from the fact that the geometry which absorbs the exhaust also governs future exhaust transport.

This tensorial step is intentionally presented as an outline rather than a finished derivation. Its extra assumptions are now explicit: a near-equilibrium local Rindler construction, a horizon-based storage rule for erased information, the load-to-stress normalization (71), and local conservation of the resulting stress-energy tensor. What the present note contributes is the bridge: the one-bit thermo-acoustic model provides an explicit microscopic instance of irreversible load generation, while the companion CCE program supplies broader conceptual context for why that load can be reinterpreted as curvature [5, 8]. On that reading, gravity is not added on top of observation; it is the large-scale bookkeeping geometry of observation itself.

### 3.6 de Sitter Noise Floor and the Deep-Field Acceleration Scale

The tensorial closure above still leaves one phenomenological question open: what fixes the scale at which the linear transport law (52) must fail? In (82), the cosmological constant enters as an integration constant. If that constant is positive, the asymptotic vacuum substrate is de Sitter rather than Minkowski, and it carries the Gibbons-Hawking temperature [6]

$$T_{\mathrm{dS}} = \frac{\hbar c}{2\pi k_B}\sqrt{\frac{\Lambda}{3}}. \tag{83}$$

Within the present information-theoretic reading, $T_{\mathrm{dS}}$ is the irreducible background noise floor of the unperturbed vacuum channel.

A local observer that reconstructs a gravitational field of magnitude $a = |\nabla\Phi|$ simultaneously defines an Unruh scale

$$T_U(a) = \frac{\hbar a}{2\pi c k_B}, \tag{84}$$

which is just (65) with the local surface gravity read as proper acceleration. The constant-$\kappa$ closure of (52) can remain self-consistent only while this local signal dominates the de Sitter background. The natural crossover to the deep-field nonlinear regime therefore occurs when

$$T_U(a_0) = T_{\mathrm{dS}}. \tag{85}$$

Equating (83) and (84) gives

$$a_0 = c^2\sqrt{\frac{\Lambda}{3}}. \tag{86}$$

The cancellation of $\hbar$, $k_B$, and $2\pi$ is noteworthy: the threshold is fixed purely by the vacuum curvature scale and the causal-temperature correspondence.

For a representative positive cosmological constant of order $\Lambda \sim 10^{-52}\,\mathrm{m}^{-2}$, Eq. (86) gives $a_0 \sim 10^{-10}\,\mathrm{m\,s}^{-2}$. That value remains within an order of magnitude of Milgrom's empirical acceleration scale inferred from galactic phenomenology, $a_{\mathrm{M}} \sim 1.2 \times 10^{-10}\,\mathrm{m\,s}^{-2}$ [11]. In the present framework, that order-of-magnitude proximity suggests that the MOND scale may mark the point at which vacuum transport of CCE exhaust first encounters the de Sitter noise floor rather than introducing an independent ad hoc constant.

The present note still stops short of a full MOND derivation. Equation (86) fixes only the crossover scale, not the detailed nonlinear constitutive law below it. A microscopic vacuum model is still required to determine the transport response that replaces the constant $\kappa$ of (52) when $a \lesssim a_0$. Nevertheless, the result constrains any such completion: its low-acceleration scale is no longer free, but is set by $\Lambda$. If the eventual nonlinear closure reproduces the standard MOND asymptotics, then at least part of the low-acceleration mass discrepancy normally attributed to dark matter would instead be re-read as vacuum backreaction in the de Sitter-noise-limited transport regime.

## 4 Conclusion

We constructed a minimal physical observer: a single reusable bit realized by a particle in a controllable double-well potential and coupled to a compressible fluid. Signal capture can be modeled reversibly, but reset cannot. Because reset erases coarse-grained information, the observer must export entropy to the bath. In the thermo-acoustic setting studied here, that load produces a steady $1/r$ thermal halo which, in a focusing medium and within the weak-field regime, bends later signals inward and enlarges the effective capture cross-section. To an external analyst, the resulting feedback can therefore look like an attractive law rather than like the observer's own maintenance cost.

The toy model is meant as a concrete instance of the broader CCE claim: once observation is treated as materially implemented, reusable records require dissipative bookkeeping, and that bookkeeping can bias the very channel through which later observations are made. In that sense, some law-like regularities need not be wholly prior to observation; they can emerge from the coupling between observer, substrate, and irreversible reset.

The gravitational extrapolation remains a key resulting claim or conjecture. Reading the vacuum as the substrate that absorbs erased information suggests a reinterpretation of $G$ as Vacuum Informational Compliance and places the breakdown of the linear weak-field branch near a de Sitter-set crossover scale, but lifting the toy model to Einstein dynamics requires additional horizon, equilibrium, and constitutive assumptions. The next step is therefore to continue to develop a sharper derivation: medium-specific transport, finite-source structure, and observers with enough internal state to track the regularities they help induce.

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