# Emergence of the Physical Laws of a Macroscopic Observer

*Peter David Fagan, School of Informatics, University of Edinburgh*

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**Abstract**

We study how apparently fundamental laws can emerge from the physical limitations of observation itself. Starting from the Conservation-Congruent Encoding perspective, we treat an observer not as an abstract witness but as a localized finite device whose records must be stored in matter and repeatedly reset. To make that bookkeeping physically explicit, we reduce the observer to a single reusable bit realized by a particle in a symmetric double-well potential and immersed in a homogeneous thermo-acoustic medium. An incoming acoustic wave packet flips that bit reversibly through Hamiltonian dynamics. Reset is then implemented by an explicit time-dependent control protocol on the double-well potential that lowers the barrier, tilts the landscape, and returns the particle to the ready well. Because that logically irreversible protocol acts while the bit remains coupled to the surrounding fluid, it dissipates at least $k_B T_0 \Delta\mathcal{H}_{cg}$ per erase into the bath, with $\Delta\mathcal{H}_{cg} \leq \ln 2$ and equality only in the maximally uncertain one-bit limit. At reset rate $\nu$, the observer therefore becomes a localized heat source of mean power $P \geq \nu k_B T_0 \Delta\mathcal{H}_{cg}$. Steady heat conduction in three dimensions then yields the exact exterior profile $\delta T(r) = P/(4\pi k r)$, and the temperature dependence of the acoustic refractive index gives $n(r) = n_0(1 + \Gamma/r)$ with $\Gamma = \frac{P}{4\pi k n_0} \left(\partial n/\partial T\right)_{T_0}$. In the focusing regime $\Gamma > 0$, the resulting bending enlarges the observer's capture cross-section and generates ray paths mathematically isomorphic to a $1/r$ attractive potential, so an external analyst restricted to the reduced event stream can mistake self-induced bath distortion for an intrinsic attractive force. The model thereby provides a concrete thermo-acoustic route by which law-like regularities can arise as projection artifacts of observation.

## 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 [5, 1]. An observer is therefore not a mathematical ghost, but a physical device whose records must be realized and maintained in matter.

The Conservation-Congruent Encoding (CCE) framework sharpens this point by defining information directly in terms of coarse-grained dynamics [3]. 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 [6, 7, 3]. 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. A question that follows is how that loss induces backreaction and impacts later observations.

This note studies the simplest setting in which that bookkeeping feeds back into what is observed. We design a toy model with a single reusable bit stored in a double-well potential and immersed in a compressible fluid. An incoming acoustic signal flips the bit reversibly under a static Hamiltonian. Reset is not left as an abstract dissipative bracket; it is implemented by an explicit time-dependent control protocol on the potential that merges the logical alternatives and returns the particle to the ready state while dissipating at least $k_B T_0 \Delta\mathcal{H}_{cg}$ per cycle into the fluid, 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 changes the observer's future capture cross-section. The central demonstration is that regularities can arise as projection artifacts of observation under the CCE framework.

## 2 The Single-Bit Capture Model

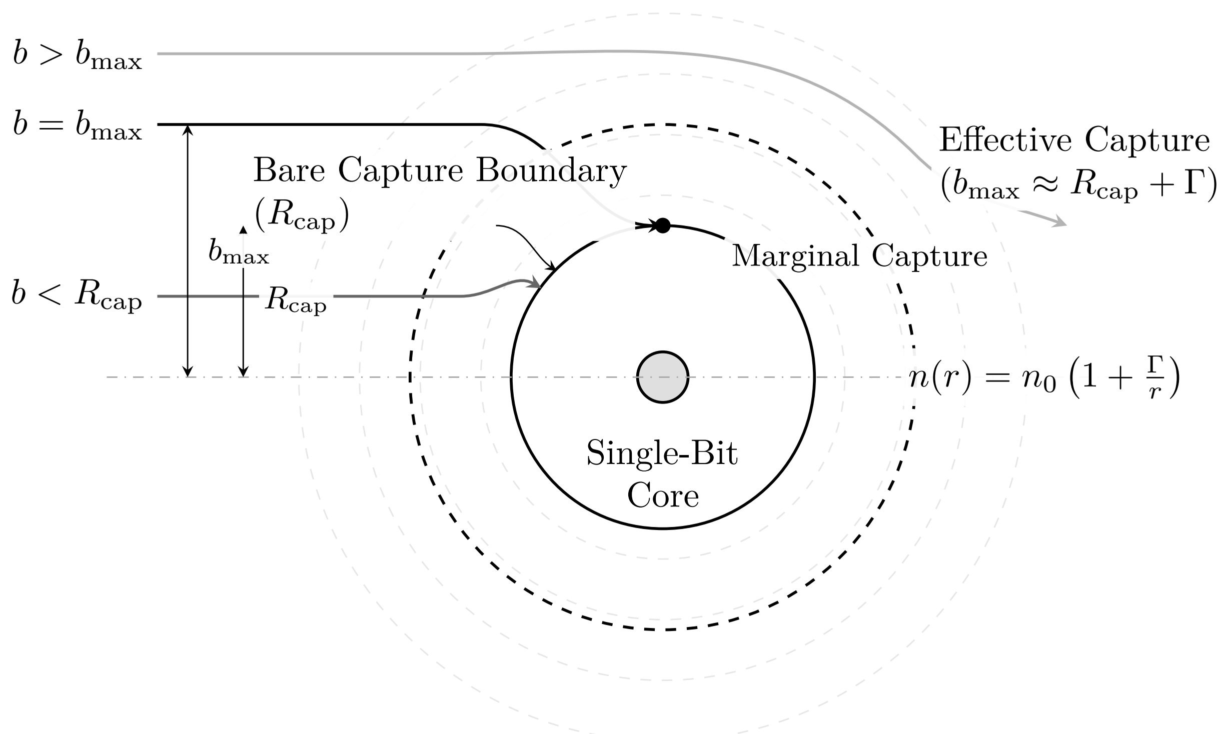


Figure 1: The effective capture cross-section of a single-bit observer. Here $R_{\text{cap}}$ denotes the effective bare capture radius of the coarse model. A time-dependent reset protocol erases the bit and dissipates heat into the surrounding fluid. In a medium with positive thermo-acoustic coefficient $(\partial n/\partial T)_{T_0} > 0$, the resulting steady $1/r$ thermal halo establishes a focusing refractive lens. Incident signals with impact parameters $R_{\text{cap}} < b \leq b_{\max}$ 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.

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

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

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

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

An incoming signal is represented by a wave packet in $\phi$ that propagates toward the observer. As it passes through $x_{\text{obs}}$, it exerts the force $\gamma\,\phi(x_{\text{obs}}, t)$ on the particle. The mechanical work transferred to the bit over the encounter is

$$W_{\text{in}} = \int_{t_i}^{t_f} \gamma\,\phi(x_{\text{obs}}, t)\,\dot{q}(t)\,dt. \tag{17}$$

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 coupling in $H_{\text{tot}}$ is pointlike, the Hamiltonian above does not by itself define a finite geometric cross-section. Rather than tracking a finite-size interaction kernel, we unify the relevant spatial scales for analytic parsimony. We therefore introduce a single coarse scale $R_{\text{cap}}$: in the absence of thermal backreaction it is the largest impact parameter for which an incoming packet still transfers sufficient work to satisfy $W_{\text{in}} \geq \Delta V$, and in the later backreaction analysis it also serves as the effective support radius of the thermal exhaust and the geometric boundary grazed by the marginal captured ray in the eikonal limit. In a refined finite-size interaction model one would derive $R_{\text{cap}}$ from the packet profile, coupling kernel, and threshold dynamics; here it is retained as the bare capture scale entering the cross-section formulas.

### 2.3 Controlled Reset and Heat Dissipation

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

$$0 \mapsto 0, \qquad 1 \mapsto 0. \tag{18}$$

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

$$W_{\text{ctrl}} = \int_{t_i}^{t_f} \frac{\partial V}{\partial t}(q(t), t)\, dt. \tag{19}$$

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

$$m\ddot{q} + \eta\dot{q} + \partial_q V(q,t) = \zeta(t), \qquad \langle \zeta(t)\zeta(t') \rangle = 2\eta k_B T_0\, \delta(t-t'), \tag{20}$$

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. The erased coarse-grained entropy per cycle is then

$$\Delta\mathcal{H}_{cg} = -p_0 \ln p_0 - p_1 \ln p_1 \leq \ln 2, \tag{21}$$

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 $T_0$,

$$\langle Q_{\text{bath}} \rangle \geq k_B T_0 \Delta\mathcal{H}_{cg}, \tag{22}$$

with equality only in the quasistatic limit [5, 2, 4]. At reset rate $\nu$, the observer therefore acts as a localized heat source of mean power

$$P := \nu \langle Q_{\text{bath}} \rangle \geq \nu k_B T_0 \Delta\mathcal{H}_{cg}. \tag{23}$$

For a maximally uncertain one-bit reset, this reduces to $P \geq \nu k_B T_0 \ln 2$. This replaces the previously postulated bath-loading rate by an explicit thermodynamic source term generated by the reset protocol itself.

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

$$\dot{A} = \{A, H_{\text{eff}}(t)\} + (A, S_{\text{eff}}). \tag{24}$$

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 Thermal 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 $R_{\text{cap}}$ denote the unified coarse capture scale introduced above. Assuming a conduction-dominated steady state over an observation time $t \gg \max(\nu^{-1}, R_{\text{cap}}^2/\chi)$, where $\chi$ is the fluid's thermal diffusivity, the repeated erasures may be replaced by a continuous mean heat source of total power $P$ supported inside $r \leq R_{\text{cap}}$. Outside that compact source region, the steady temperature perturbation

$$\delta T(r) = T(r) - T_0 \tag{25}$$

obeys the sourceless heat equation

$$\nabla^2 \delta T = 0, \qquad r > R_{\text{cap}}, \tag{26}$$

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

$$-4\pi r^2 k \, \frac{d\delta T}{dr} = P, \qquad r > R_{\text{cap}}, \tag{27}$$

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

$$\delta T(r) = \frac{P}{4\pi k r}, \qquad r > R_{\text{cap}}. \tag{28}$$

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

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

$$n(r) = \frac{c_0}{c(T(r))}, \qquad c_0 := c(T_0), \qquad n_0 := n(T_0) = 1, \tag{29}$$

and linearizing around the ambient state gives

$$n(r) = n_0 \left[1 + \beta_T \, \delta T(r)\right] + \mathcal{O}\!\left(\delta T(r)^2\right), \qquad \beta_T := \frac{1}{n_0}\left(\frac{\partial n}{\partial T}\right)_{T_0}. \tag{30}$$

Substituting the exact thermal profile yields

$$n(r) = n_0 \left(1 + \frac{\Gamma}{r}\right), \qquad \Gamma = \beta_T \frac{P}{4\pi k}. \tag{31}$$

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 $\beta_T = -1/(2T_0) < 0$, so the focusing version of the model requires a different thermo-acoustic medium. By contrast, simple organic liquids such as ethanol or benzene can exhibit $\partial c/\partial T < 0$ over ordinary laboratory conditions, which implies $\beta_T > 0$ and supplies a concrete focusing medium for the toy model.

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 $R_{\rm cap}$ and the scale over which $n(r)$ varies. In that eikonal limit, propagation may be treated in terms of rays obeying Fermat's principle.

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

$$\Delta\theta(b) \approx -\frac{1}{n_0}\int_{-\infty}^{\infty} \frac{\partial n}{\partial b}\,dx = \Gamma \int_{-\infty}^{\infty} \frac{b\,dx}{(x^2+b^2)^{3/2}} = \frac{2\Gamma}{b}. \tag{32}$$

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

$$\sigma_0 = \pi R_{\rm cap}^2. \tag{33}$$

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

$$L = n(r)\, r \sin\psi, \tag{34}$$

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 capture surface at $r = R_{\rm cap}$ with $\sin\psi = 1$. In the focusing regime $\Gamma > 0$, conservation of $L$ therefore gives the exact threshold

$$b_{\max} = \frac{n(R_{\rm cap})}{n_0} R_{\rm cap} = R_{\rm cap}\left(1 + \frac{\Gamma}{R_{\rm cap}}\right) = R_{\rm cap} + \Gamma, \tag{35}$$

and hence

$$\sigma_{\rm eff} \approx \pi\left(R_{\rm cap} + \Gamma\right)^2 = \pi R_{\rm cap}^2 + 2\pi R_{\rm cap}\Gamma + \mathcal{O}(\Gamma^2). \tag{36}$$

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

$$\Gamma \ge \beta_T \frac{\nu k_B T_0 \Delta\mathcal{H}_{cg}}{4\pi k}, \qquad (\beta_T > 0), \tag{37}$$

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

$$\sigma_{\rm eff} - \sigma_0 \gtrsim \frac{\beta_T R_{\rm cap} \nu k_B T_0 \Delta\mathcal{H}_{cg}}{2k} \tag{38}$$

in the weak-lensing regime. To close the observational feedback loop, we take the observer to run an unconditional reset clock tuned to the expected hit rate. If the ambient incident flux is $J$, we therefore set the mean clock rate by $\nu \approx J\sigma_{\rm eff}$, so the lower-bound estimate becomes

$$\sigma_{\rm eff} \gtrsim \sigma_0 + \Lambda J \sigma_{\rm eff}, \qquad \Lambda := \frac{\beta_T R_{\rm cap} k_B T_0 \Delta\mathcal{H}_{cg}}{2k}, \tag{39}$$

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

$$\sigma_{\rm eff} \approx \frac{\sigma_0}{1 - \Lambda J}. \tag{40}$$

For a maximally uncertain one-bit reset, $\Delta\mathcal{H}_{cg} = \ln 2$ and the earlier $\ln 2$ scaling is recovered. The observer's own erasure cost therefore increases its future capture probability through a concrete thermo-acoustic channel, with an incipient runaway capture regime as $\Lambda J \to 1^-$. An external analyst who sees only the enhanced hit rate can 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 mimics a fundamental force. 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 $p(r) \approx p_0\,(1 - U(r)/(m_{\text{test}} v_0^2))$ reveals that the ray paths are identical to the orbits 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 an effective inward radial acceleration of $a_r = -v_0^2\Gamma/r^2$. The same thermodynamic exhaust that enlarges $\sigma_{\text{eff}}$ thus also projects the exact orbital geometry of an inverse-square attractive field.

# 3 Conclusion

This note has recast the toy model in concrete physical terms. The observer is a localized finite system with exactly one reusable bit, stored as a particle in a controllable double-well potential and immersed in a compressible fluid. Capture is a fully Hamiltonian process: an incoming acoustic wave packet transfers enough work to push the particle across the static barrier $\Delta V$. Reset is then executed by an explicit time-dependent control protocol that lowers the barrier, tilts the potential, and restores the ready state while the bit remains coupled to the fluid. Because that protocol erases coarse-grained information, it dissipates at least $k_B T_0 \Delta\mathcal{H}_{cg}$ per cycle into the bath, with $\Delta\mathcal{H}_{cg} \le \ln 2$ and equality only in the maximally uncertain one-bit limit, so at reset rate $\nu$ the observer acts as a localized heat source of mean power

$$P \ge \nu k_B T_0 \Delta\mathcal{H}_{cg}. \tag{41}$$

When the observer instead runs an unconditional reset clock tuned to the expected hit rate from an ambient incident flux $J$, this rate is not exogenous but closes self-consistently as $\nu \approx J\sigma_{\text{eff}}$.

Steady Fourier conduction then gives the exact exterior temperature profile

$$\delta T(r) = \frac{P}{4\pi k r}, \tag{42}$$

and linear thermo-acoustic response converts that halo into the refractive law

$$n(r) = n_0\left(1 + \frac{\Gamma}{r}\right), \qquad \Gamma = \beta_T \frac{P}{4\pi k}. \tag{43}$$

In the focusing regime $\beta_T > 0$, the capture threshold rises to $b_{\max} = R_{\text{cap}} + \Gamma$ and the effective cross-section becomes $\sigma_{\text{eff}} \approx \pi R_{\text{cap}}^2 + 2\pi R_{\text{cap}}\Gamma$. Furthermore, through the optical-mechanical analogy, the geometry of Fermat ray paths maps exactly to classical orbits in an attractive potential, yielding an apparent radial acceleration $a_r = -v_0^2\Gamma/r^2$. The thermal parameter $v_0^2\Gamma$ thus plays the exact role that an external Newtonian description would assign to a gravitational parameter $GM$. What looks like an external attractive source is therefore, in this toy model, the observer's own irreversible bookkeeping fed back through a concrete thermo-acoustic bath.

That thermo-acoustic chain is the single-channel realization of the more general CCE bookkeeping. In the present model the dissipative load is thermal, so the emitted power obeys the generalized CCE lower bound

$$P \geq \mathcal{J}_{c,\min} = (\mathcal{F}_c \alpha_c)\, \Delta\mathcal{H}_{cg}. \tag{44}$$

Here the reset throughput is $\mathcal{F}_c = \nu$ and the channel conversion factor is $\alpha_c = k_B T_0$, so the abstract CCE current becomes the thermal lower bound $\mathcal{J}_{c,\min} = \nu k_B T_0 \Delta\mathcal{H}_{cg}$. In the maximally uncertain one-bit limit, $\Delta\mathcal{H}_{cg} = \ln 2$ and the bound reduces to $P \geq \nu k_B T_0 \ln 2$. The toy model therefore does not bypass the CCE formalism; it instantiates it in a particularly transparent thermo-acoustic setting.

The transport step also abstracts cleanly. In this note the erased entropy is carried away by heat, and steady Fourier conduction generates the scalar load $\delta T(r) = P/(4\pi k r)$. But the same exterior geometry is not uniquely thermal. Any localized observer that repeatedly exports a conserved burden into a three-dimensional medium may produce an analogous diffusive halo when the surrounding substrate is governed by a comparable linear transport law. In a different substrate, similar mathematics might arise for electrochemical charge, spin angular momentum, or other conserved loads, but each case would require its own transport model.

The backreaction step is likewise universal even though its constitutive coefficient is medium-specific. In the thermo-acoustic realization that susceptibility is

$$\beta_T = \frac{1}{n_0} \left( \frac{\partial n}{\partial T} \right)_{T_0}, \tag{45}$$

which converts the observer's thermal exhaust into a refractive bias on future signals. In another substrate, the analogous response coefficient could instead modify magnetic permeability in a superconducting environment, ionic conductivity in a biological membrane, or some other transport property that controls what later signals reach the observer. These examples are illustrative rather than derived here; each would require a medium-specific constitutive analysis. The microscopic mechanism changes from medium to medium, but the CCE lesson does not: once irreversible reset loads the surrounding substrate, the substrate can feed that load back into the observer's future event stream as a regular bias.

The significance of the model is therefore conceptual rather than phenomenological. It demonstrates a physical realization of the emergence mechanism in one concrete example and then shows, through the CCE mapping, why the same logic may scale beyond thermo-acoustics. A fuller treatment could derive $\beta_T$ and $k$ from a specific fluid model, resolve the finite-size source structure inside $r \leq R_{\text{cap}}$, and go beyond the weak-lensing estimate used here. It would also be useful to study when additional internal state—for example, a clock or short event buffer—lets the observer itself infer the regularities that in the present note are reconstructed only by an external analyst. Even in stripped-down form, however, the Single-Bit Capture Model makes the central point sharply: once observation is treated as materially implemented, conserved-current exhaust and substrate backreaction can make law-like behavior emerge from the architecture of observation itself.

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