

Thermal Avalanches Drive Logarithmic Creep in Disordered Media

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Disordered materials under an imposed forcing can display creep and aging effects, accompanied by intermittent, spatially heterogeneous dynamics. We propose a testable microscopic description of these phenomena, based on the notion that as the system ages, the density of local barriers that enable relaxation displays a slowly evolving gap. As a result, the relaxation dynamics is dominated by the activation of the lowest, extremal tail of the distribution. This framework predicts logarithmic creep, as well as correlated bursts of slow activated rearrangements, or “thermal avalanches,” whose size grows logarithmically with their duration. The time interval between events within avalanches obeys a universal power-law distribution, with a cutoff that is simply proportional to the age of the system. We show that these predictions hold both in numerical models of amorphous solids, as well as in experiments with thin crumpled sheets. This analysis suggests that the heterogeneous dynamics occurring during logarithmic creep is related to other phenomena, including dynamical heterogeneities characterizing the glass transition.

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I. INTRODUCTION

Disordered systems in even the simplest situations can give rise to rich emergent physics. Consider the behavior of a thin sheet of plastic crumpled into a ball. This seemingly mundane system exhibits a wide range of far-from-equilibrium behaviors, including slow relaxations, creep and aging effects [1,2], intermittent response accompanied by crackling sounds [3] and avalanche statistics [4], as well as strong memory effects [1,2,5,6]. Under a change of external load, the crumpled sheet compacts logarithmically over many decades in time, exhibiting “creep” flow. Such logarithmic creep is ubiquitous to disordered solids and has been observed over timescales ranging from seconds to days (or even years), including in crumpled sheets [1–4], disordered electronic systems [7–9], frictional contact strength [10–12], polycrystalline ice [13], granular compaction [14–18], the motion of pinned elastic interfaces [19,20], and in amorphous solids [21–26]. The experimental accessibility of crumpled sheets offers a view of the

microscopic processes underlying logarithmic creep. It was recently shown that the logarithmic compaction of crumpled sheets under load advances via ever-slowing, scale-free cascades of mesoscopic instabilities [4]. However, a complete microscopic description of the observed intermittent and slow behavior and the resulting logarithmic aging is still missing.

Scale-free distributions of rapid activity bursts, commonly known as crackling noise [27], are commonly observed in disordered driven systems [28,29]. These phenomena are mostly studied theoretically at zero temperature via mean-field approximations, scaling and renormalization group arguments [30,31], or by considering the self-organized criticality of cellular automata [32–34]. At finite temperatures, the situation is more complex. “Thermal avalanches” of slow activated events have been argued for in that case, but their nature is debated [35–40]. Thermal avalanches are anomalously slow: The time interval between events constituting them can be orders of magnitude longer than for mechanical avalanches, where it is controlled by the speed of sound. Instead, thermal avalanches are governed both by interactions and thermal noise: Each instability lowers the energy barriers of subsequent events, facilitating their thermal activation at a possibly much later time. References [36,37] predicted that their size grows logarithmically in time and are responsible for dynamical heterogeneities in supercooled liquids near their glass transition

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[41–46] and in the stationary creep flows of pinned interfaces [36,47]. However, empirical support for this anomalous temporal dynamics is limited to molecular dynamics of model liquids near their glass transition [37,48] or immediately after a quench [40], and experimental validation is missing. Moreover, this approach considered stationary conditions, and its relation to aging phenomena is unclear.

In this paper, we extend these views to obtain a microscopic description of logarithmic creep. Our approach, which focuses on the evolution of the energy barrier distribution in the material, naturally captures the logarithmic slowdown observed in creep as well as the observed statistics of thermal avalanches. It also makes novel predictions on the distribution of time intervals between instabilities within avalanches. This approach ties together perspectives from dynamical heterogeneities in supercooled liquids [37,45,48], activated avalanches in depinning models [36,39], and classical exhaustion theories where creep is caused by a depletion of weak regions in the materials over time, initially developed for crystalline plasticity [21].

In a second step, we successfully test these predictions in experiments with crumpled thin sheets [4] and in numerical simulations of model amorphous solids under an applied load [49,50]. Overall, our work leads to a novel framework to describe aging occurring during creep, which connects microscopic and macroscopic phenomena in a variety of driven systems.

II. A MESOSCOPIC APPROACH

We model a disordered system of linear size L and dimension d in a mesoscopic manner, as being composed of $N = L^d$ blocks, which satisfy the following properties:

- (1) Each block i is characterized by a stability index x_i , quantifying the additional forcing that must be applied to the block to rearrange it. For example, for amorphous solids under an applied shear, x_i is often considered to represent the difference between the shear component of the local stress tensor and that of the local yield stress.
- (2) A stability index x implies that the block displays an energy barrier for relaxation that vanishes when $x = 0$, and therefore, for small x , $E = \epsilon_0 x^\alpha$. The constant ϵ_0 can be taken to be unity without loss of generality for a proper choice of the scale of stability values, whereas α is an exponent that depends on the smoothness of the local energy barrier ($\alpha = 3/2$ if the energy landscape is smooth and therefore exhibits a saddle-node bifurcation [51–53], and $\alpha = 2$ if it presents cusps [54]).
- (3) At finite temperature T (in units where the Boltzmann constant is $k_B = 1$), each stable site can activate and rearrange with an Arrhenius rate $\lambda(x) = \exp[-E/T]/\tau_0$. Note that τ_0 is a microscopic

timescale that henceforth defines our units of time, such that $\tau_0 = 1$.

- (4) Rearranging blocks are coupled with some interaction kernel $\delta x_j = G(\vec{r}_{ij})$ describing the change in stability at block j after a relaxation event in block i . Note that $G(\vec{r}_{ij})$ can be local and only couple nearby blocks, or it can be long ranged. For amorphous solids, this kernel represents the spatial stress redistribution, which generically decays as $G(\vec{r}_{ij}) \sim 1/r^d$ and can change sign depending on the direction considered [55].

This mesoscopic approach has a long history in various fields, including the depinning transition [31,39,53], the glass transition [37,45,46,56,57], the plasticity of amorphous solids [55,58–60], and, more recently, in crumpled thin sheets [4,61].

Below, we develop a theoretical description of logarithmic creep flow under an external load, based on this common mesoscopic view. Crucial to our results will be the microscopic interplay between thermal activation and interactions [37,45]. We expect the main predictions of our theory to apply to systems satisfying assumptions 1–4. Our description is agnostic to the specific value of α . The shape of the interaction kernel (e.g., short range vs long range) can affect the universality class and the value of critical exponents but not the overall scenario we propose. Finally, our theoretical framework may hold if thermal fluctuations are replaced by mechanical noise (e.g., vibrations in the experimental room) as long as the latter is not too broadly distributed.

III. EXISTING ARGUMENTS FOR STATIONARY FLOWS

Before describing creep dynamics, let us first review results that describe steady-state flow using the above mesoscopic perspective. At vanishing temperature $T = 0^+$, the dynamics in this viewpoint is approximately “extremal”: The weakest site (i.e., the smallest x_i value, which we denote $x_{\min}$) relaxes first [19]. Extremal dynamics has been studied extensively in the context of self-organized criticality [32,33,62]. It has been shown that, for stationary conditions, as $N \rightarrow \infty$, the distribution of stability $P(x) = \sum_i \delta(x - x_i)/N$ develops a gap at a finite stability value x_g . However, below the gap, there is a subextensive population of stable sites [36,62]. In the thermodynamic limit, the smallest stability value $x_{\min}$ is distributed as

$$P(x_{\min}) \sim (x_g - x_{\min})^{-\beta}, \quad (1)$$

where β is some exponent, which, in some cases, can be derived analytically [62]. As is the case for all exponents introduced in this work, β can depend on the shape of the propagator $G(\vec{r}_{ij})$. Note that the value of β plays a negligible role for the aging dynamics we derive below, as it only enters

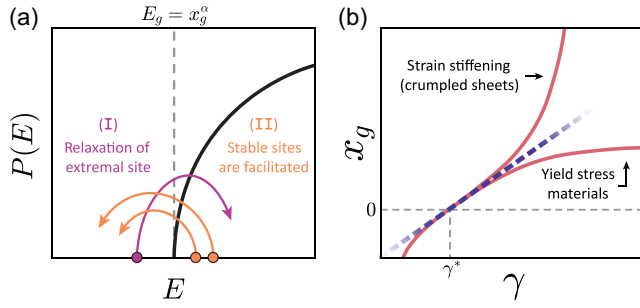


FIG. 1. Theoretical scenario. (a) Dynamical facilitation in the density of local barriers, which characterizes the energy landscape. The extremal site with lowest energy (typically close to E_g) is activated by thermal fluctuations. It facilitates stable sites to lower energies. These sites are then activated quickly, altogether forming a thermal avalanche. (b) Sketch for the stress gap x_g as a function of strain γ at small temperatures and fixed external stress σ . For $T = 0^+$, extremal dynamics is a good approximation, and it leads to a stability gap x_g that is zero at $\gamma = \gamma^*$ and initially grows linearly with strain. The linear approximation of $x_g(\gamma)$ close to γ^* (dashed line), which lies at the base of our theory, can be applied both to strain stiffening materials and to yield stress materials.

in logarithmic corrections to the asymptotic behaviors we study.

At finite temperature, it was argued in Ref. [36] that extremal dynamics remains a good approximation below some length scale $\xi \sim T^{-\nu}$ characterizing dynamical heterogeneities. For smaller systems $L < \xi$ as considered in what follows, thermal avalanches can be defined as cascades of thermally activated events driven by facilitation [37]. In essence, each local instability can lower the barriers of neighboring stable sites, facilitating their thermal activation [62]. This process is illustrated in Fig. 1(a). For any arbitrary threshold time t_0 , sequences of S events separated by time intervals shorter than t_0 are power-law distributed with

$$P(S) \sim S^{-\tau} g(S/S_c), \quad (2)$$

where $g(z)$ is a cutoff function that decays sharply from 1 to 0 around $z \approx 1$. Reference [37] showed that the cutoff S_c satisfies

$$S_c \sim [T \ln(t_{\text{sta}}/t_0)]^{-\nu d_f}, \quad (3)$$

where d_f is the fractal dimension. Here, $t_{\text{sta}} \equiv \exp[E_g/T]$ is the characteristic relaxation time of the system in the stationary state. In essence, thermal avalanches are anomalously slow, and their size grows logarithmically with the cutoff time interval t_0 that defines them. Because the overall duration of an avalanche is dominated by the longest intervals between events constituting it, this duration is also of order t_0 . Equation (3) was never tested empirically, and here we test this relation using experiments on crumpled sheets.

IV. EXTENSION TO AGING DYNAMICS

We now consider a system where a load is imposed at $t = 0$. In the limit of small temperatures, the dynamics will be approximately extremal as least stable sites will yield much faster than others. At long times, the system must tend towards the stationary description, with a gap in the distribution of stabilities [36,62]. We thus expect a gap x_g to rapidly form; yet, its value must slowly evolve to adapt to the applied load. This view is consistent with experiment and modeling of crumpled sheets, for which the system's relaxation slows down due to a slow but steady increase in the system's lowest (local) energy barrier [4]. Thus, the dynamics will be similar to the stationary case described above: Microscopic relaxation is characterized by long waiting times of order $\exp[x_g^\alpha/T]$ at which thermal avalanches are triggered. Here, however, x_g evolves in time, thus accounting for the aging-related slowdown.

A. LOGARITHMIC CREEP

To characterize this evolution, we denote by γ the quantity conjugate to the applied load (corresponding to the strain in amorphous solids or crumpled sheets). Here, $x_g(\gamma)$ must be a smooth function in the extremal dynamics approximation. Denoting γ^* the strain at which mechanical stability is reached, i.e., $x_g(\gamma^*) = 0$, we must then have $x_g(\gamma) \approx C_0(\gamma - \gamma^*)$ in the neighborhood of γ^* , which is illustrated in Fig. 1(b). The strain rate $\dot{\gamma}$ is inversely proportional to the long waiting time between thermal avalanches set by x_g and thus must follow at the leading order:

$$\dot{\gamma} \sim \exp[-(\gamma - \gamma^*)^\alpha C_0^\alpha/T]. \quad (4)$$

In the Appendix, B we show that this equation leads approximately to

$$\gamma - \gamma^* \sim [T \log(t)]^{1/\alpha}, \quad (5)$$

$$\dot{\gamma} \sim \frac{1}{t}, \quad (6)$$

where logarithmic prefactors are neglected in the last expression. This analysis thus recovers the logarithmic response observed as such systems age. From Eq. (5), we also obtain for the gap energy controlling the relaxation time:

$$E_g \sim T \log(t). \quad (7)$$

This derivation is reminiscent of phenomenological strain hardening models used to account for logarithmic aging [1,63] and also of “exhaustion theories” of creep [21,64]. These models assume that the system's evolution is dictated by a single energy barrier that grows in proportion to the overall strain. Here, it is derived from first principles

as a result of extremal activation of the distribution of barriers near the gap.

B. THERMAL AVALANCHES

The increasing stability gap leads to a growing timescale $t_g \equiv \exp[-E_g(\gamma)/T]$. Since avalanches are subextensive, the effect of a single avalanche on the gap (an intensive quantity) must be vanishing in the thermodynamic limit. Thus, as far as a single avalanche is concerned, the gap is essentially constant in time. The results on avalanche statistics in the stationary case must then apply: These avalanches will also exhibit a cutoff S_c as in Eq. (3), with the evolving t_g playing the role of t_{sta} so that

$$\langle S_c \rangle \sim [T \ln(t_g/t_0)]^{-\nu d_f}, \quad (8)$$

with prefactors independent of t_0 that slowly grow in time before saturating, as detailed in the Appendix. Compared to the stationary case, the difference lies in the timescale t_g , which is not constant and instead grows linearly as the system ages.

C. WAITING TIME DISTRIBUTION

During thermal creep, instabilities follow universal temporal statistics. Assuming extremal dynamics, so that creep is dominated by the waiting time of the weakest site [19], with $P(E_{\min}) \sim (E_g - E_{\min})^{-\beta}$ from Eq. (1), then $P(\Delta t/t_g) = \int P(\Delta t/t_g|U)P(U)dU$, where $U = E_g - E_{\min}$. Since the waiting time Δt distribution for a site with barrier $E_{\min}$ is exponential, we have that $P(\Delta t/t_g|U) = \lambda e^{-\lambda \Delta t/t_g}$ for Arrhenius rate $\lambda = \exp[U/T]$. For $\beta = 0$, an exact integration is possible, leading to

$$P(\Delta t/t_g) \approx \left(\frac{\Delta t}{t_g}\right)^{-1} \exp\left[-\frac{\Delta t}{t_g}\right], \quad (9)$$

which holds for $\Delta t \gg 1$, and where a logarithmic prefactor in t_g is neglected, as reported in the Appendix. This formulation demonstrates that the distribution of $\Delta t/t_g$ essentially does not depend on t_g , as we will confirm empirically below. When $\beta > 0$, there are additional logarithmic corrections (see the Appendix for a derivation).

V. EXPERIMENTS AND SIMULATIONS

To verify our theory, we consider the experimental setup introduced in Ref. [4]. Thin sheets of Mylar, 50 × 50 cm across and 8 μm thick, are crumpled manually several times and placed under an external load of $M = 400$ g. Note that M is well below the load applied during the manual crumpling, such that it does not generate new creases or induce significant plastic flow [65]. Instead, it perturbs the existing, precrumpled structure [4]. We monitor the height of the crumpled ball $h(t)$.

Over time, a steady logarithmic compaction $h(t) = h_0 - b \ln(t)$ is observed, typically reaching relative compaction of order $[h_0 - h(t)]/h_0 \approx 5\%$. As the sheet compacts, it emits crackling noise [3,66], intermittent acoustic pulses that are picked up by a nearby microphone and assigned a time stamp t_i . These originate from mesoscopic snap-through instabilities in the sheet [5,6]. Reference [4] has established that these instabilities govern the logarithmic compaction of crumpled sheets, while the contribution of plastic flow in individual creases [67] in between instabilities can be neglected.

We complement the experiments with a minimal mesoscopic thermal elastoplastic model (EPM) of amorphous plasticity, as described and implemented in Refs. [49,50]. Similar implementations of this thermal EPM have been used to model the avalanches and rheology of steady-state flow in amorphous solids near the critical flow stress at finite temperature [50] as well as transient creep flow [22,23]. Here, we use EPMs to study dynamical heterogeneities during creep. The model itself consists of a grid of L^2 elastically coupled sites, each of which represents a coarse-grained block of material that can potentially rearrange. Each site is associated with a local stress threshold $\sigma_{\text{th},i}$ that, when exceeded, allows for plastic rearrangement in the site. The annealed disorder in this system comes from sampling the initial and rejuvenated $\sigma_{\text{th},i}$ thresholds from a Weibull distribution with shape parameter $k = 2$. The localized rearrangement that occurs when $\sigma_i > \sigma_{\text{th},i}$ allows the site to locally strain (by approximately $L^{-d}\delta\gamma$) into a more favorable configuration, thereby reducing the local stress σ_i . This process induces a long-range stress redistribution throughout the system, stabilizing some sites and destabilizing others, with an Eshelby-like kernel computed using finite elements, which is asymptotically $G(\vec{r}_{ij}) \sim \cos(4\theta)/r^d$, where θ is the angle between $\vec{r}_{ij}$ and the direction of applied external shear. We allow for the thermal activation of otherwise mechanically stable sites, (i.e., with $x_i = \sigma_{\text{th},i} - \sigma > 0$) with an Arrhenius rate $\lambda(x) = \exp[-x^\alpha/T]$ (having set $k_B = 1$). We choose $\alpha = 3/2$, as predicted from catastrophe theory for smooth particle interactions [51,68] (though the choice of α only affects logarithmic corrections), and $T = 10^{-3}$ as a rescaled temperature for which extremal dynamics is an excellent approximation. In this low-temperature limit, the EPM exhibits self-organized criticality [60,69–71], as it is a distant cousin of sandpile and Bak-Sneppen-type models [32,62] and avoids finite-size-induced continuous thermal flow (as occurs for $T > T^*(L)$ [36,50]). We load our EPM simulations at a fixed stress $\Sigma = 0.5$, well below the critical yield stress $\Sigma_c \approx 0.73$ for our model.

Despite the distinct differences in the rearrangement mechanisms, the interactions, and the fate under large deformation (namely, strain stiffening vs flowing), both experiments and simulations are captured by our theoretical analysis. They exhibit logarithmic thermal creep

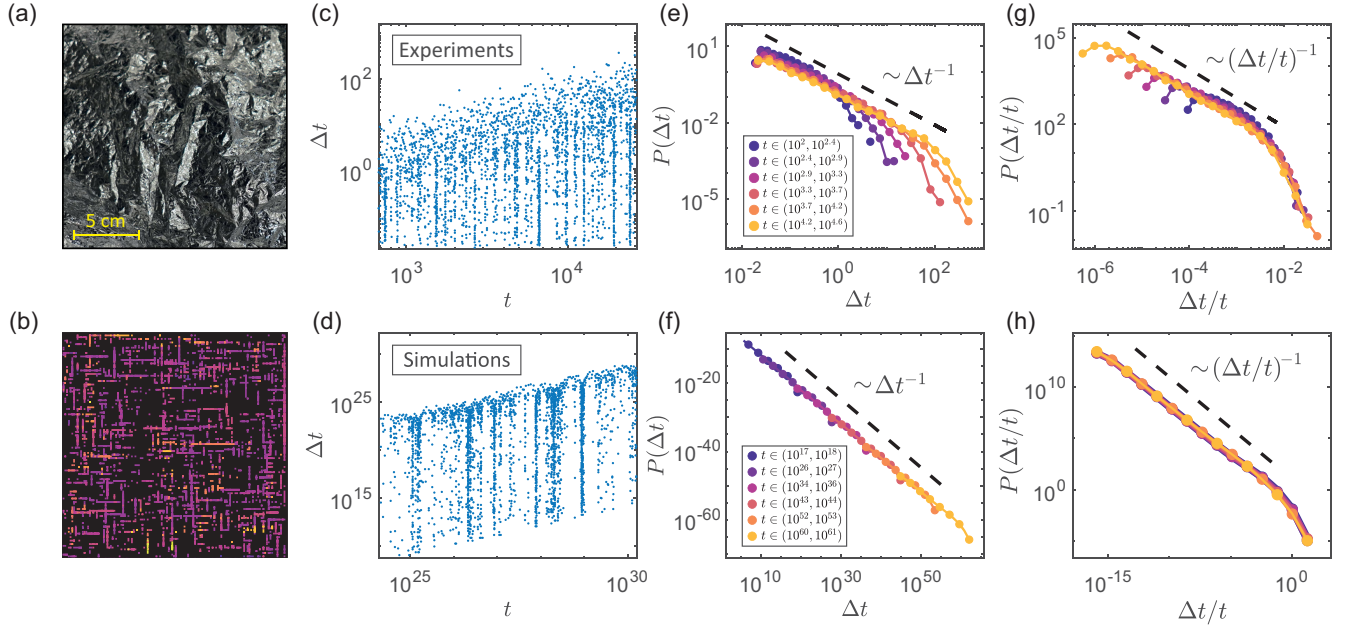


FIG. 2. Temporal statistics of instabilities. Instabilities during logarithmic creep in experiments (top) and simulations with $L = 128$ (bottom): (a) unfolded crumpled sheet of Mylar; (b) visualization of the EPM's activity over a strain interval of $\Delta\gamma \approx 0.4$. Instability locations are colored by the log of their occurrence time, revealing many linelike avalanches. (c),(d) Waiting times Δt between instabilities as a function of time. Each point represents a single instability. (e),(f) Waiting time distributions $P(\Delta t)$ in different time windows going from purple to yellow. All distributions follow $P(\Delta t) \sim \Delta t^{-1}$. (g),(h) Normalizing by t , the distributions $P(\Delta t/t)$ collapse to a single curve for the same time windows. For the experimental results, time is measured in seconds.

via avalanches of instabilities, corresponding to an opening gap in the distribution $P(x)$. Most prominently, this dynamics is evident when plotting the waiting time between instabilities $\Delta t_i = t_i - t_{i-1}$ vs their occurrence time t_i [4], as shown in Figs. 2(c) and 2(d). During creep, the typical waiting time between instabilities grows, yet we observe thermal avalanches of instabilities that involve long timescales. These avalanches are evidenced by the vertical streaks in Figs. 2(c) and 2(d). Figure 2(b) shows an example of the spatiotemporal dynamics of activity in the EPM, where the characteristic “linelike” avalanche implies a fractal dimension $d_f \approx 1$. From the envelope of the stabilities and waiting times between instabilities, we can access the gaps x_g and t_g (see the Appendix for detailed protocol).

Note that the extremely small temperature T that drives the EPM results in an astronomically long transient creep phase [72]. Along with an event-based algorithm [49,50], this choice allows us to probe logarithmic creep for over 50 time decades (timescales only accessible in simulations), providing a stringent test of our scaling predictions.

VI. AGREEMENT WITH THEORY

We begin by considering waiting time distributions $P(\Delta t)$. We divide the experiments and simulations into logarithmically spaced intervals, such that, in each bin, the

cutoff t_g is approximately constant. As predicted by extremal dynamics of Arrhenius crossings, $P(\Delta t)$ in each interval follows the universal power law of Eq. (9), as shown in Figs. 2(e) and 2(f). Indeed, the logarithmic corrections to Eq. (9) can be neglected. The cutoff grows linearly with time $t_g \sim t$, and the distributions over different intervals collapse when normalizing by time and calculating $P(\Delta t/t)$, as shown in Figs. 2(g) and 2(h).

Next, we consider the energy landscape and its evolution during creep. The simulations allow us to access the full barrier distribution $P(E)$ at a given time. As predicted, for low E , the distribution develops a gap over time, as shown in Fig. 3(a). Note that the usual power-law behavior $P(E) \sim (E - E_g)^{\tilde{\theta}}$ is observed above the gap, as expected in the presence of long-range interactions with varying sign [69,73,74]. Importantly, below this gap, the distribution is nonzero (inset), though it is fluctuating and subextensive with system size [62]. Creep is dominated by these barriers below the gap E_g . Indeed, isolating only the density of instabilities that were triggered, E_{fail} , reveals a wide distribution peaking at $E_g(t)$, as shown in Fig. 3(b). As detailed in the Appendix, we use statistical methods to extract the gap magnitude E_g and confirm Eq. (7): $E_g \sim T \ln(t)$. We find that the barriers are power-law distributed relative to the gap $P(E) \sim (E_g - E)^{-\beta}$ with $\beta \approx 0.4$, as shown in Fig. 3(c). This result is consistent

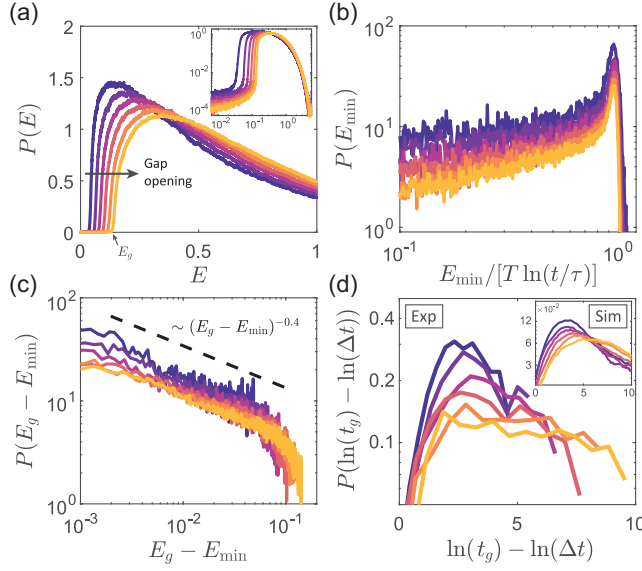


FIG. 3. Evolution of the energy landscape. (a) Full energy barrier distributions $P(E)$ across different ages in simulations. With time, the distribution develops a gap E_g . Color coding is consistent with the time windows in Fig. 2. The inset shows the same plot in log-log scale. (b) Barrier distributions for failing sites only, $P(E_{\text{fail}})$, as a function of E_{fail} , normalized by the gap $E_g \sim T \ln(t)$. The collapsed curves reveal rich activity below the gap E_g at different times. (c) Barrier distribution of failing sites relative to the gap, exhibiting a power law $P(E) \sim (E_g - E)^{-\beta}$ with $\beta \approx 0.4$. (d) Barriers estimated from waiting times as $E/T = \ln(\Delta t)$, leading to similar results in the experiments and the simulations (inset).

with recent simulations of thermally activated, disordered interfaces, where a power-law distribution of extremal energy barriers develops below a fixed energy gap [36].

In experiments, the distribution of failing sites can only be accessed indirectly, by approximating $E_{\text{fail}} \sim \ln(\Delta t)$ and $E_g \sim \ln(t_g) \sim \ln(t)$. In contrast to the actual barrier energy distribution, the distribution relative to the gap $P(\ln(t_g) - \ln(\Delta t))$ exhibits a peak and a sharp decay close to zero, as shown in Fig. 3(d). This effect originates from fluctuations in the barrier crossing time Δt : Although a barrier of height E predicts a mean waiting time of $\tau \exp[E/T]$, the waiting time distribution is exponential, meaning that longer times are possible. Since fluctuations are on the order of the mean crossing time, $\ln(t_g)$ is, in fact, an overestimation of the gap when compared to extracting it directly. Indeed, the same distribution $P(\ln(t_g) - \ln(\Delta t))$ exhibits a peak in the simulations as well, and it resembles the experimental results (Fig. 3(d) inset). This effect hinders the ability to extract the logarithmic correction β from the experimental data, but direct correspondence between numerics and experiments supports our approach.

We now examine the distribution of thermal avalanches. The magnitude of each avalanche S is calculated as the number of consecutive events with $\Delta t_i < t_0$. Whenever t_0 is

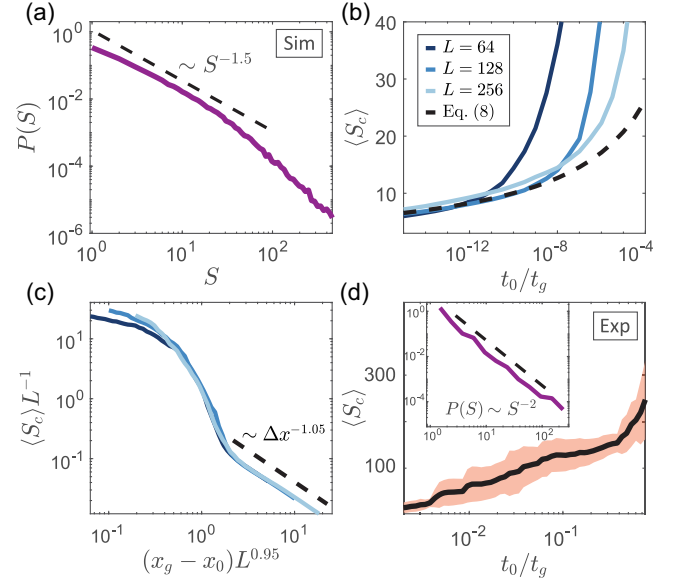


FIG. 4. Thermal avalanches. Each avalanche is calculated as the sequences of events with $\Delta t_i < t_0 < t_g$, while the gap is $t_g = Ct$. (a) Avalanche distribution in simulations with $L = 256$, and for fixed $t_0/t_g = 10^{-4}$, exhibiting a power-law decay $P(S) \sim S^{-\tau}$, with $\tau = 1.5 \pm 0.1$. The exponent τ is obtained via maximum likelihood estimation. (b) Avalanche cutoff $\langle S_c \rangle$ as a function of t_0/t_g in simulations, and a fit to Eq. (8) with $\nu d_f \approx 1.05$. As system size is increased, the fit holds closer to the gap and for larger t_0/t_g . (c) Finite-size scaling of the avalanche cutoff, plotted by the distance to the stability gap $\Delta x = x_g - x_0$. This reveals a minimum stability scale of about $L^{-0.95}$, above which our theory holds. (d) Thermal avalanches in experiments. The inset shows the avalanche distribution for $t_0/t_g = 10^{-1}$. The distribution exhibits a power law with $\tau = 2 \pm 0.2$. The main panel shows the avalanche cutoff $\langle S_c \rangle$ as a function of t_0/t_g (solid curve), which agrees qualitatively with the theoretical prediction. The shaded region marks the fluctuations between experiments.

exceeded, a new avalanche begins. Crucially, t_0 is not constant but instead grows linearly with time $t_0 = At$, such that the distance to the gap t_0/t_g remains constant during aging. Indeed, for $t_0/t_g < 1$, we find that the avalanche distributions are well captured by Eq. (2) and exhibit scale-free distributions with $\tau \approx 2 \pm 0.2$ in experiments and $\tau \approx 1.5 \pm 0.1$ in simulations (as determined by maximum-likelihood estimation [75]). This case is shown in Figs. 4(a) and 4(d). The different exponents we find in experiments and simulations could arise from the functional form of the propagator $G(\vec{r}_{ij})$ in EPM ($d = 2$) and the crumpled sheets, a fractal object ($2 < d < 3$ [76]). Nevertheless, we note that detecting events experimentally is limited to $\Delta t > 0.02$ s, potentially leading to an underestimation of magnitudes S and therefore of τ .

We can now test the prediction of Eq. (8). We vary t_0/t_g systematically and compute the thermal avalanche distribution for each value. The distribution cutoff $\langle S_c \rangle$ is obtained from the moment ratio $\langle S_c \rangle \sim \langle S^{k+1} \rangle / \langle S^k \rangle$ [77].

As a function of t_0/t_g , it exhibits the predicted slow logarithmic growth until diverging close to $t_0/t_g \rightarrow 1$, which is well captured by Eq. (8) in the numerical model, as shown in Fig. 4(b). Here, we find $\nu d_f \approx 1.05$, consistently with previous studies [69].

This fit succeeds only up to a system-size-dependent ratio $t_0/t_g < r_{\text{coal}}(L)$, which marks the onset of coalescence between different thermal avalanches. In other words, the ability to separate avalanches is limited by finite-size fluctuations, and as $t_0 \rightarrow t_g$, these fluctuations lead to the coalescence of independent cascades. Figure 4(c) shows how larger systems can probe the dynamics closer to the gap, extending the $s_c \sim (x_g - x_0)^{-1/\sigma}$ scaling to a minimum stability scale $x_g - x_0 \sim L^{-0.95}$. Accordingly, very large avalanches are formed exclusively through coalescence, and they belong to a distinct percolation universality class that is an artifact of our measurement strategy in finite systems. A similar transition occurs in neuronal avalanches when coalescence becomes dominant [70]. In the Appendix, we address this issue with a full finite-size scaling analysis.

For the experiments, the growth of the cutoff $\langle S_c \rangle$ qualitatively matches the expected behavior. It exhibits a slow growth and diverges when $t_0/t_g \rightarrow 1$. However, finite-size fluctuations, along with our temporal resolution, do not allow a precise estimation of νd_f or a fit to Eq. (8).

VII. DISCUSSION

We have proposed a microscopic framework for the aging and creep flows of various disordered materials. In this description, a key quantity is the distribution of local barriers, in particular, its external tail. We have described how this density adiabatically ages in time by opening up a gap. This view naturally explains ubiquitous observations, including the macroscopic logarithmic response to an imposed perturbation and the presence of an intermittent microscopic dynamics. Concerning the latter, we predicted that it occurs by bursts of slow activated events, whose stability is smaller than that of the gap. Quantitatively, the size of these bursts is logarithmic in their duration. The distribution of time laps between events collapses when rescaled by the system age and follows an inverse power law. The vast differences between the systems we studied, crumpled sheets and amorphous solids, suggest that our description may apply broadly. It would be particularly interesting to test our views in granular systems. In that case, local stabilities presumably reflect weakening inter-grain contacts [78,79], which may present local energy barriers coupled elastically. If so, we expect thermal avalanches to be present in that system as well, lasting much longer than mechanical avalanches propagating at the speed of sound.

Central to our description is the appearance of a time-dependent gap in the distribution of energy barriers and the recognition that creep proceeds through the activation of

sites below this gap. Thus, the distribution of activated barriers has a sharp cutoff, $P(E_{\min} > E_g) = 0$. The weakest energy barrier is essentially never above E_g , owing to the repopulation of sites below the gap by facilitation. This feature contrasts with renewal models such as Bouchaud's trap model [80], where the exponential tail of a *fixed* barrier distribution $\exp[-E_{\min}/T_g]$ gives rise to a power-law distribution of waiting times, and aging. Alternative descriptions of logarithmic creep consider a wide range of independent relaxation modes [81], leading to log-Poisson statistics [3]. While the exponential cutoff of Eq. (9) is reminiscent of log-Poisson dynamics [82], the approximately inverse power law reveals that the dynamics is far more complex and highly correlated.

Our scenario incorporates the notion that dynamical heterogeneities are induced by logarithmically growing thermal avalanches. This view has received numerical support in stationary conditions in systems as diverse as pinned interfaces [36] or molecular liquids near their glass transition [48]. Our work provides a quantitative, empirical test for the dynamics of thermal avalanches, in addition to extending this approach to nonstationary conditions. Such heterogeneous dynamics may thus be generic to a broad class of disordered materials and phenomena and could be revealed, for example, by examining the distribution of waiting times $P(\Delta t)$. Future studies could explore the complex spatial and temporal correlations within thermal avalanches, which may reveal deeper insights into their collective organization [47].

Our framework offers new perspectives to study complex aging phenomena, including nonmonotonic aging [2,83,84], memory [2,85,86], and rejuvenation effects [87]. Which observable in these materials retains the history of past perturbations imposed on them? A natural hypothesis that stems from our approach is that memory is encoded in the extremal tail of the distribution of local barriers. Finally, a central question in both material science and geophysics is what events can nucleate rupture. Sudden rupture events can occur, for example, when very stable amorphous materials are prepared and then loaded [88,89] or when velocity instabilities are induced by velocity-weakening effects as in faults or frictional interfaces [90]. An interesting hypothesis to investigate in the future is the notion that thermal avalanches may act as seeds that nucleate shear bands or cracklike propagating ruptures.

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APPENDIX A: CUTOFF OF AVALANCHE SIZES

In simulations, we consider avalanches defined as sequences of events for which the minimal stability $x_{\min} \leq x_0$. The cutoff for the avalanche distribution is known to scale as [62]

$$s_c \sim (x_g - x_0)^{-\nu_{df}}. \quad (\text{A1})$$

For small enough temperatures in a finite system, an equivalent formulation allows us to define avalanches strictly in terms of experimentally accessible waiting times: A t_0 avalanche proceeds without any quiescent periods longer than t_0 . We have

$$\log(t_g/t_0) = \exp[(x_g^\alpha - x_0^\alpha)/T]. \quad (\text{A2})$$

Using that at first order $x_g^\alpha - x_0^\alpha \approx \alpha x_g^{\alpha-1}(x_g - x_0)$ leads to

$$\log(t_g/t_0) = \exp[(x_g^\alpha - x_0^\alpha)/T] \approx \exp[\alpha x_g^{\alpha-1}(x_g - x_0)/T]. \quad (\text{A3})$$

Rearranging, and using Eq. (A1) to eliminate $x_g - x_0$, we obtain

$$s_c \sim [T \log(t_g/t_0)]^{-\nu_{df} x_g^{\nu_{df}(\alpha-1)}}. \quad (\text{A4})$$

In a stationary state or when approaching it, x_g is approximately constant. At early time during creep, we instead have $x_g \approx [T \log(t_g)]^{1/\alpha}$, indicating that avalanches may tend to grow (logarithmically) in time if $\alpha > 1$.

APPENDIX B: GAP EVOLUTION SOLUTION

The equation for the growth of strain justified in the main text follows Eq. (4).

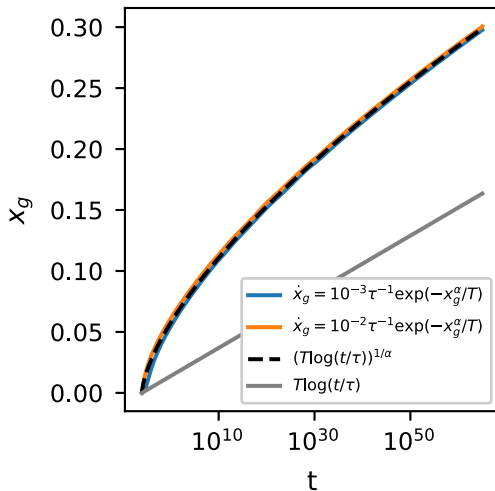


FIG. 5. Numerical integration of Eq. (4) yielding a solution $x_g \approx [T \log(t)]^{1/\alpha}$.

We are not aware of a closed-form solution of that equation when $\alpha \neq 1$. It is easy to check that the form of Eq. (5) approximately solves it, except for subdominant logarithmic terms. The numerical solution of Eq. (4) is shown in Fig. 5, confirming the validity of this approximation.

APPENDIX C: WAITING TIME DISTRIBUTION

Recall that the time interval distribution, normalized by t_g , follows $P(\Delta t/t_g) = \int \lambda e^{-\lambda \Delta t/t_g} P(U) dU$, where $U = E_g - E_{\min}$, $P(U) = [(1-\beta)/(E_g^{1-\beta})] U^{-\beta}$, and $\lambda = \exp[U/T]$. Thus, $P(\Delta t/t_g) = \int e^{U/T - \exp[U/T] \Delta t/t_g} U^{-\beta} dU$. For $\beta = 0$, an exact integration is possible, yielding

$$p\left(\frac{\Delta t}{t_g}\right) = \frac{T}{E_g} \left(\frac{\Delta t}{t_g}\right)^{-1} \left[\exp\left(-\left(\frac{\Delta t}{t_g}\right)\right) - \exp(-\Delta t) \right]. \quad (\text{C1})$$

The second exponential is only relevant when $\Delta t \lesssim 1$, but it is necessary to ensure that the distribution is normalizable. The overall normalization is maintained by a factor of T/E_g , and noting that $T/E_g = \log(t_g)$, this gives the simpler expression

$$p\left(\frac{\Delta t}{t_g}\right) \approx \frac{1}{\log(t_g)} \left(\frac{\Delta t}{t_g}\right)^{-1} \exp\left[-\left(\frac{\Delta t}{t_g}\right)\right]. \quad (\text{C2})$$

For our experimental and numerical data, we cannot measure $\Delta t/t_g$ below some $\bar{t}_{\min} > 1$. Thus, the overall normalization does not depend on $1/\log(t_g)$ but instead on $\log(1/\bar{t}_{\min})$.

For the case $\beta \in [0, 1)$, we resort to a Laplace approximation, where the integral is estimated at the value where the exponential contribution is maximum. To lowest order, this occurs at $U^* = T \log(t_g/\Delta t)$. In the limit of $\Delta t/t_g \ll 1$, we find approximately

$$p(\Delta t/t_g) \sim \frac{(\frac{\Delta t}{t_g})^{-1} (T \ln(\frac{t_g}{\Delta t}))^{1-\beta}}{\sqrt{(\ln(\frac{t_g}{\Delta t}))^2 - \beta}}. \quad (\text{C3})$$

This expression is valid in the limit where $T\beta \ll 1$, $\Delta t/t_g \ll 1$, and $\Delta t \gg 1$.

APPENDIX D: GAP ESTIMATION IN EXPERIMENTS AND SIMULATIONS

To precisely assess the stability gap x_g in simulations, we consider the time series of failing sites (t_i, x_i) and calculate its envelope. Namely, at time t , the gap is

$$x_g(t) = \max_{t_i \leq t} \{x_i\}. \quad (\text{D1})$$

This property is averaged over 60 realizations, resulting in a smooth, linearly growing stability gap. The gap energy in Fig. 3 is simply $E_g = x_g^\alpha$.

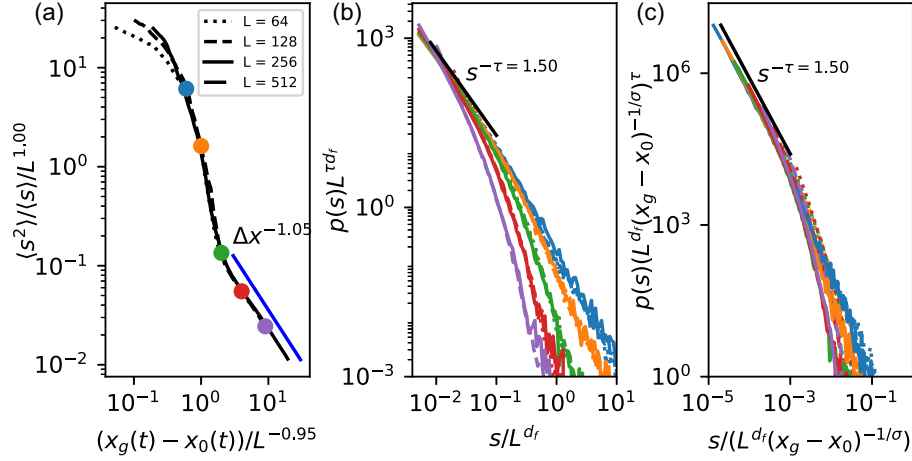


FIG. 6. (a) Growth of $s_c(x_0)$ as a function of δx_g , rescaled to effect a finite-size scaling collapse. (b) Avalanche distributions, evaluated for different system sizes and at different values of $\delta x_g/L^{-a}$ [indicated by the colored circles of panel (a)], also rescaled to effect a collapse in system size. (c) Same avalanche distributions, rescaled by $s_c \sim \delta x_g^{-1/\sigma}$ to effect a collapse.

Similarly, for the waiting times in Fig. 3(d) and 4, we compute the envelope

$$t_g(t) = \max_{t_i \leq t} \{\Delta t_i\}. \quad (\text{D2})$$

In experiments, there is a larger variability between different realizations. Namely, the prefactor C for the linear gap growth $t_g = Ct$ is not constant across all experiments. Thus, instead, we calculate the gap by Eq. (D2) for each realization separately and perform a linear fit to obtain a smoothly growing gap.

An alternative method to assess the growth of the gap is by the waiting time distribution cutoff, which can be accessed from the moment ratio $t_c \sim \langle \Delta t^3 \rangle / \langle \Delta t^2 \rangle$. Calculating the moment ratio at different ages also yields a linearly growing gap. However, the prefactor for the moment ratio depends on the waiting time probability density function, including the logarithmic correction β . Because of this sensitivity, we choose the first method, which is robust regardless of $P(\Delta t)$.

APPENDIX E: FINITE-SIZE SCALING OF THE GAP

When the stability gap is measured using the envelope [Eq. (D1)] in a finite system, there will be fluctuations of order $\Delta x_g(L) \sim L^{-a}$ that make it impossible to accurately delineate thermal avalanches below this scale. In other words, in an infinite system, x_0 avalanches exhibit a cutoff scaling as $s_c \sim (x_g - x_0)^{-1/\sigma} = \delta x_g^{-1/\sigma}$. In a finite system, this scaling holds only for $\delta x_g > \Delta x_g(L) \sim L^{-0.95}$, as can be seen in Fig. 6(a). For $\delta x_g < \Delta x_g(L)$, the x_0 avalanches begin to merge together into much larger events. These events do not represent causal cascading activity from a single activation of an initiating site near x_0 but rather the merging of several consecutive x_0 avalanches near the

fluctuating $x_g(t, L)$ gap. The rapid growth of s_c for $x_g - x_0 < \Delta x_g(L)$ reflects a transition to this new merging mode of growth. Accordingly, the avalanche distribution [Fig. 6(b)] exhibits a two-power-law decay. The first power law relates to the causal avalanches of sites below the gap while the second captures the effect of randomly merging cascades together. Similar distributions are encountered in neural systems when independently initiated power-law cascades (directed percolation) merge together with a distinct universality class (undirected percolation) [70]. The scaling of s_c with $\Delta x_g^{-1/\sigma}$ only applies to the initial cascades, which is why the collapse in Fig. 6(c) only applies to systems with $\Delta x_g > 0$.

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