Intermittent collective dynamics of domain walls in the creep regime

Matías Pablo Grassi,1 Alejandro B. Kolton,1,2 Vincent Jeudy,3 Alexandra Mougin,3 Sebastian Bustingorry,4 and Javier Curiale1,4

1Instituto Balseiro, Universidad Nacional de Cuyo - CNEA, Av. Bustillo 9500, 8400 S. C. de Bariloche, Río Negro, Argentina.
2CONICET, Centro Atómico Bariloche, 8400 San Carlos de Bariloche, Río Negro, Argentina.
3Laboratoire de Physique des Solides, CNRS, Univ. Paris-Sud, Université Paris-Saclay, 91405 Orsay, France.
4Instituto de Nanociencia y Nanotecnología, CNEA–CONICET, Centro Atómico Bariloche, 8400 San Carlos de Bariloche, Río Negro, Argentina.

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We study the ultra slow domain wall motion in ferromagnetic thin films driven by a weak magnetic field. Using time resolved magneto-optical Kerr effect microscopy, we access to the statistics of the intermittent thermally activated domain wall jumps between deep metastable states. Our observations are consistent with the existence of creep avalanches: roughly independent clusters with broad size and ignition waiting-time distributions, each one composed by a large number of spatio-temporally correlated thermally activated elementary events. Moreover, we evidence that the large scale geometry of domain walls is better described by depinning rather than equilibrium universal exponents.

I. INTRODUCTION

Domain walls (DW) in thin ferromagnetic films have become a paradigmatic system1–6 to learn about the universal interplay between disorder, elasticity and thermal fluctuations in driven interfaces. Such physics is relevant for a large variety of experimental systems7–9, and for potential applications as DW are building blocks for proposed magnetic storage devices10. The caveat is that even an arbitrarily weak disorder has a rather dramatic effect on the DW dynamics, notably the occurrence of a depinning threshold11–14. Below the threshold, DW are pinned at zero temperature and they present a thermally activated glassy behavior called the creep regime at finite temperature. A better understanding of the impact of disorder in the low velocity regimes is thus fundamental for a comprehensive study of DW dynamics, and of disordered elastic interfaces in general.

Most of the experimental studies on weakly driven DW motion, including very recent ones15,16 focused on the universal features of the steady DW mean velocity vs the field $H$ and temperature $T$, but not in its spatio-temporal fluctuations. Such kind of study has been mostly performed close to the depinning threshold where the fluctuations are dominated by large deterministic collective events. For example, avalanche size distribution and its universal properties has been discussed in the context of Barkhausen noise17, contact lines of liquids18, crack propagation19, and even in reaction fronts in disordered flows20 and active cell migration21. Well below the depinning threshold, the phenomenology of avalanches have remained much less clear22. Recently however, theoretical studies of ultra-slow creep motion23–25 have unveiled rather unexpected and non trivial spatio-temporal patterns, whose elementary events strongly differ from those encountered close to the depinning threshold. Therefore, tackling experimentally a detailed statistical study of magnetization reversal events is particularly interesting.

The numerical simulations reported in Ref. 26 show that creep motion of a one dimensional interface model proceeds via a sequence of elementary events (EE) of fluctuating sizes. These EE are the minimal thermally activated jumps that make the DW overcome energy barriers and irreversibly advance under the applied field $H$. The size statistics of EE display broad distributions, with a characteristic lateral size cut-off $L_{opt} \sim H^{-3/4}$ and a characteristic area size $S_{opt} \sim L_{opt}^{3/4}$. These results confirm the existence of an optimal “thermal nucleus”, as proposed in the pioneer creep theories11,12. Since energy barriers for DW motion scale as $U_{opt} \sim L_{opt}^{1/3} \sim H^{-1/4}$ (Ref. 17), Arrhenius activation of these nuclei leads to the celebrated creep-law $\ln v \sim -H^{-1/4}/T$ for the mean velocity $v$ at which the DW move under the action of a small magnetic field $H$. The EE are not normally distributed in size and are not independent as traditionally assumed. On one hand, below $S_{opt}$, EE areas are power-law distributed as $P_{EE}(S) \equiv S^{-\tau_E}G(S/S_{opt})$, with $\tau_E$ a characteristic exponent and $G(x)$ a rapidly decaying function for $x > 1$. On the other hand, EE tend to cluster in space and time forming larger cluster events (CE). These CE are similar to the so called “creep avalanches” suggested by functional renormalization group calculations in Ref. 17 and experimentally noticed in Ref. 15. Such composite objects are, unlike EE, weakly correlated and have a much broader distribution of areas, $P_{CE}(S) \sim S^{-\tau_{CE}}$ with $\tau_{CE}$ a universal exponent. These interesting predictions were not yet evidenced experimentally nor confirmed by other theoretical approaches.

In this work we test the above scenario by a statistical analysis of the ultra slow time evolution of magnetization reversal in ferromagnetic Pt/Co/Pt thin films. For different time windows of duration $\Delta t$, we determine the size ($S$) distribution $P_{EE}(S) \equiv P_{EE}(S; \Delta t, T, H)$ of the ob-
served consecutive compact magnetization reversal area that we call “window-event” (WE). This procedure permits us to relate WE with EE and CE and to show that the features displayed by $P_{gg}(S)$ are consistent with the picture summarized above of rare localized EE acting as epicenters of large CE or “creep avalanches”, each made of a large number of spatio-temporally correlated EE. Furthermore, our analysis of the intermittent collective DW motion allows to characterize the statistics of waiting times between epicenter EE, thus going beyond the “geometric” predictions of Ref [16].

II. METHODS

Experiments were mainly performed on a Pt(4.5 nm)/Co(0.7 nm)/Pt(3.5 nm) thin ferromagnetic film with perpendicular magnetic anisotropy. A polar magneto-optical Kerr effect (PMOKE) microscope was used to image magnetic domains. In order to characterize the DW dynamics, starting with a seed magnetic domain, used to image magnetic domains. In order to characterize magneto-optical Kerr effect (PMOKE) microscope was applied perpendicular to the film plane to favour the growth of the initial domain. The DW velocity was then computed following a standard differential protocol. After identifying the creep regime in the $H-T$ plane by fitting the creep-law $\ln v \sim -H^{-1/4}/T$, we fix $T$ to two possible values, room temperature and 50 °C, and choose $H = 46.1$ Oe and $H = 24.2$ Oe respectively, such that $v \sim 1$ nm s$^{-1}$ in each case. We then analyze the magnetization reversal events at each temperature, for a total applied field time $t = 27000$ s. Since the characteristic areas of EE are expected to scale as $S_{opt} \sim H^{-5/4}$, and the energy barriers for nucleation as $U_{opt} \sim H^{-1/4}$, choosing fields deep in the creep regime allows us to maximize, in principle, our spatial and temporal sensitivity to intrinsic collective events. For these fields we indeed observe a clear intermittent (i.e. not smooth) growth.

To characterize it statistically, during the long-time magnetic field pulse we stroboscopically observe the growth at intervals $\Delta t$, such that $t \gg \Delta t$. The duration $\Delta t$ is much larger than the acquisition time of each image, and much smaller than the pulse time $t$ so to collect a large number of events. This allows them to compute their area distribution, $P_{gg}(S)$, for different $\Delta t$ and $T$. Although we mainly report results for one region of a specific sample, we have also performed less detailed but similar measurements in other regions of the same sample and also in a different material and checked robustness of our results. We discard WE touching any border of the region of interest in order to not underestimate their area and make a proper comparison with theoretical predictions. We have checked that this protocol does not affect the tails of $P_{gg}(S)$ for the time windows $\Delta t$ used. We refer the reader to appendix A for further details on our experimental setup and protocols. Magnetization reversal events were previously obtained in irradiated Pt/Co/Pt samples[13], identifying between 30 and 50 events depending on field values. In the present work we were able to obtain thousands of WE, thus allowing a more precise statistical description, amenable to comparison with the universal theoretical predictions.

III. RESULTS

A. Domain wall motion within the creep regime

The obtained field dependence of domain wall velocity for the analyzed sample is presented in Fig. 1(a). The figure shows the evolution of the velocity as a function of the magnetic field over eight orders of magnitude. Within the creep regime, thermal activation over a field dependent energy barrier leads to a stretched exponential increase of the velocity, given by [11,12,17]

$$v = v_0 \exp \left[ -\frac{T_d}{T} \left( \frac{H}{H_d} \right)^{-\mu} \right],$$

where $v_0$ is a temperature dependent velocity, $T$ the temperature, $k_B T_d$ a typical energy scale coming from the competition between elasticity and disorder ($k_B$ being the Boltzmann constant), $H_d$ the depinning field and $\mu = 1/4$ the universal creep exponent. As shown in Fig. 1(b), a straight line with a negative slope in a plot of $\ln v$ against $H^{-1/4}$ confirms that the measured velocities are within the creep regime, and in addition that the system belongs to the universality classes of one dimensional elastic systems displacing in a two dimensional media, with a random-bond type of disorder and short-range elasticity. The fit to the creep formula of Eq. (1) for the two temperatures we analyzed are

$$\ln[v \text{ m}^{-1}] = -128(1)(\text{Oe})^{1/4}H^{-1/4} + 27.6(3),$$

$$= -100(2)(\text{Oe})^{1/4}H^{-1/4} + 24.4(5)$$

at $T = RT$ [Eq. (2)] and $T = 50^\circ C$ [Eq. (3)]. This data, the experimental estimates for the depinning field $H_d$, the depinning temperature $T_d$ and key characteristic scales are reported in Table I.

| $T$ [K] | 293 | 323 |
| $H$ [Oe] | 46.1 | 24.2 |
| $T_d$ [K] | 7142 | 8369 |
| $H_d$ [Oe] | 760 | 650 |
| $H_{d/4}T_d^{-1} \equiv \left( \frac{H_d}{T_d} \right)^{1/4}$ | 128 | 100 |
| $T/T_d$ | 0.04 | 0.06 |
| $\left( \frac{T}{T_d} \right) \left( \frac{H_d}{T_d} \right)^{1/4}$ | 0.02 | 0.02 |

TABLE I. Characteristic depinning values for the studied temperature and field values: $H_d$ is the depinning field, $T_d$ is the depinning temperature, $H_{d/4}T_d^{-1} \equiv \left( \frac{H_d}{T_d} \right)^{1/4}$ is the slope of the creep plot [see Eqs. (2) and (3)], and $T/T_d$ and $\left( \frac{T}{T_d} \right) \left( \frac{H_d}{T_d} \right)^{1/4}$ are related to the distribution of waiting times as discussed in Sec. IIIE.
The creep velocity law: $v \sim H^{-1/4}$, where $H$ is the magnetic field. The red point corresponds to a single long pulse of duration $t = 27000$ s at a small field of $H = 46.1$ Oe and room temperature (RT). The total reversed area over this long pulse is indicated in (c) and corresponds to $v = 1.7 \times 10^{-9}$ m/s. During this long pulse PMOKE images were taken every $t_0 = 15$ s, allowing to identify $N_{\text{WE}}(t_0, t) = 1151$ magnetization reversal events or “window events” (WE), highlighted over the image. The color scale corresponds to the time at which each WE was observed.

With the aim of pursuing the characterization of small magnetization reversal events responsible of the creep motion of elastic systems, one should consider that the typical area size $S_{\text{opt}}$ of the “optimal thermal nuclei” responsible for the velocity of Eq. (1) dramatically increases when the magnetic field decreases, as $S_{\text{opt}} \propto H^{-\nu_{\text{opt}}(1+\zeta_{\text{opt}})}$, with $\nu_{\text{opt}}$ and $\zeta_{\text{opt}}$ positive universal exponents. Since velocity follows an stretched exponential dependence with $S_{\text{opt}}$, decreasing the magnetic field implies to perform very long time experiments. Therefore, after nucleation of a single domain, a small magnetic field ($H = 46.1$ Oe) is applied during a single long time pulse ($t = 27000$ s, i.e. 7.5 hours), reaching a velocity $v = 1.7 \times 10^{-9}$ m/s. The velocity-field data thus obtained is indicated as a red point in Fig. 1(b), and the differential image shown in Fig. 1(c) corresponds to the full displacement of the domain wall under these conditions. In order to identify magnetically reversed regions (WE), during the whole long time pulse experiment, the total reversed area (indicated in Fig. 1(c)) is fragmented into many small spatially compact regions obtained from the subtraction of consecutive images taken after $t_0$. The number of WE, is $N_{\text{WE}}(t_0, t) = 1151$ and are highlighted with a color code over the image of the reversed area in Fig. 1(c). Due to the characteristics of the used PMOKE microscope and the image analysis, the smallest detectable displacement of the domain wall correspond to events close to $0.3 \mu m^2$ (25 pixels).

### B. Event areas

![Event areas](image)

In Fig. 2 we show typical WE sequences, for four different values of $\Delta t$, from a $15$ s to $120$ s. We can appreciate that, for a given growth, each $\Delta t$ induces a particular partition of the total reversed area of the sequence. At large $\Delta t$ the coalescence of several smaller WE corresponding to smaller $\Delta t$ becomes evident.

In Fig. 3(a),(b) we compare size distributions $P_{\text{WE}}(S)$, from $\Delta t = 15$ to $180$ s at room temperature $T = RT$ and a field $H = 46.1$ Oe, and from $\Delta t = 20$ to $160$ s at $T = 50$ °C and a field $H = 24.2$ Oe, respectively. The first remarkable feature of all these distributions is their broadness, which can be roughly described by $P_{\text{WE}}(S) = S^{-\tau_{\text{WE}}}G_{\text{WE}}(S/S_{\text{WE}})$, where $\tau_{\text{WE}}$ is an effective power-law exponent and $S_{\text{WE}}$ the cut-off value such that the function $G_{\text{WE}}(x)$ is constant for small $x$ and decays faster than a power-law for $x \gtrsim 1$. Quantitatively similar size-distributions were observed in different regions of the same sample and also in other kind of magnetic films (see appendix A).
FIG. 3. WE area distributions for increasing window times ∆t (as indicated) at RT and H = 46.1 Oe (a) and at T = 50 °C and H = 24.2 Oe (b). In both cases v ~ 1 mm s⁻¹. At small S we compare the initial decay of P_EE(S) with S⁻¹τ_EE, with τ_EE ≈ 1.11, where τ_E calls for depinning avalanches. The collapse scaling shows that the data of (a) and (b) displays a large size cut-off scaling S_EE ∼ (∆t/τ_t)¹/², with τ_t an H and T dependent characteristic time. (d) Effective power-law exponents τ_EE for P_EE(S) vs ∆t/τ_t.

Both τ_EE and the large-size cut-off S_EE depend on ∆t. As can be appreciated in Fig. 3(a)-(b) S_EE increases with ∆t, more specifically S_EE ∼ (∆t/τ_t)¹/². The fair collapse of P(S)∆t¹/² vs S/∆t¹/² shown in Fig. 3(c) confirms this dependence. Here, t* ≡ t*(T, H) is a characteristic time. Concomitantly, in Fig. 3(d) we show that τ_EE ≈ 1 for the smallest ∆t/τ_t for the whole data of Figs 3(a),(b). Note also that the same t* that describes the S_EE (T, H) dependence allows to build a master curve for τ_EE vs. ∆t/τ_t. For the characteristic times t* we find t*_{50°C} ≈ 1 s at T = 50 °C, H = 24.2 Oe and t*_{RT} ≈ t*_{50°C}/3 at T = RT, H = 46.1 Oe. Therefore S_EE ∼ (∆t)¹/² µm² s⁻¹/² in the first case, and S_EE ∼ (3∆t)¹/² µm² s⁻¹/² in the second one.

Since EE of Ref. 16 are power-law distributed with an exponent τ_EE ≈ 1.17 it is tempting to directly compare small ∆t WE, which are also typically small, to EE. A rough estimate for the Pt/Co/Pt films we study shows that the largest EE are of the order of S_{opt} = 10⁻³(H_d/H)¹/²5 µm², where H_d is the depinning field. Since H_d ≈ 637 Oe, and our lowest field is H = 46 Oe, we get that S_{opt} ∼ 10⁻⁵ µm², which is clearly small below our PMOKE resolution of roughly 0.5 µm² (25 pixels). We thus conclude that our detected WE can not be single EE, but the sum of a large number of them. Namely, if in a time window ∆t we have N_{EE} such events, of sizes s₁, s₂, ..., s_{N_{EE}}, compactly grouped in a WE, its random area is S_{EE} ≈ ∑_{i=1}^{N_{EE}} s_i. The statistics of S_{EE} thus directly relates to the statistics of EE random sizes s_i contributing to the same WE and of their ∆t dependent and fluctuating number N_{EE}.

Given the small area of the EE compared to our detected WE, a pure statistical analysis is convenient. If the EE were considered independent and accumulating at a well defined rate on each WE, by virtue of the central limit theorem we would naively expect P_{EE}(S) to develop an approximate gaussian shape around N_{EE}. P_{EE}(S) shows no tendency to approximate a normal nor even a peaked distribution however: it is broad, even for ∆t in the minutes time scale. To interpret this it is worth recalling that the central limit theorem tell us that S_{EE} ≈ ∑_{i=1}^{N_{EE}} s_i should converge to a Gaussian distribution if N_{EE} is large enough and the s_i have finite variance and short-ranged correlation. The EE have finite variance and, although they appear to be spatially correlated, there is no evidence of correlation between their areas We hence interpret that N_{EE} must be a strongly fluctuating quantity for all the ∆t analysed. Indeed, we experimentally observe for a fixed ∆t well defined bursts of magnetic activity, with S_{EE} > 0.3 µm², coexisting with WE in the resolution edge S_{EE} > 0.3 µm², at the same H and T. Since any PMOKE resolved area S_{EE} > 0.3 µm² has a large number of EE we arrive to the first important observation of our paper: EE are strongly clustered spatio-temporally.

C. Domain Wall roughness

The results of the previous section are consistent with the EE clustering predicted for simple domain wall models [10,20]. To go beyond, since EE are too small to be experimentally resolved, one is immediately tempted to compare our experimentally resolved WE with the predicted CE. Indeed, unlike EE, CE are not expected to be strongly correlated as we also observe for WE. Moreover, the predicted value for τ_CE ≈ 1.11 is only slightly above τ_EE ≈ 1 observed in Fig. 3(d) for the smallest ∆t. To argue that WE may indeed approach the single intrinsic CE in the small ∆t limit, we start by noting that the same scaling of zero temperature depinning avalanches, S_i ∼ L^1/z_{d}, is also expected for CE at finite temperature. In Fig. 4(a) we analyse for T = RT the approximately obtuse shapes of WE by plotting the areas S_i of each WE sampled from a long sequence, versus their corresponding lateral size L_i, defined as the major axis length of the reversed blobs. A crossover is observed at S ≈ 2 µm² below which we observe a S_i ∼ L_i¹/z_{d} scaling [21]. The two main candidates depining universality classes that are consistent with the observed creep law in v ∼ H⁻¹/₄/T are the 1d quenched-Edwards-Wilkinson (qEW), and the 1d quenched-Kardar-Parisi-
Zhang (qKPZ). The first predicts $\zeta_d \approx 1.25^{22}$ while the second $\zeta_d \approx 0.6^{23,25}$ Only the qEW value is in good quantitative agreement with Fig. 3(a), in the small size WE limit$^{23}$. In addition Fig. 3(b) is quantitatively consistent with the relation $\tau_{CE} = 2 - 2/(1 + \zeta_d) \approx 1.11$ predicted for qEW. To investigate this issue in Fig. 3(b) we computed the squared width $W^2(L) = \frac{u^2(x) - u_L(x)^2}{2}$ from different small segments of size $L$ extracted from typical DW configurations, where $u_L(x)$ is the DW displacement measured with respect to the untilted segment (see Appendix B details). The scaling $W^2 \sim L^{2\zeta_d}$ is consistent with the qEW depinning roughness exponent $\zeta_d = 1.25$ and thus with Figs. 3(d) and 3(a). We then arrive to the second important observation of our paper: WE approach single CE in the small $\Delta t$ limit and we find experimental evidence that the DW roughness and CE statistical properties are better described by depinning rather than equilibrium exponents as theoretically predicted in Refs. $^{17-27}$ and $^{28}$. A smaller roughness exponent $0.69 \pm 0.07$ was observed in extended DW in the same material in the pioneering work by Lemerle et al. $^{11}$ and interpreted to be the equilibrium exponent $\zeta_{eq} = 2/3$. Such interpretation implies an observation scale below $L_{opt}$ $^{24}$. However, from Ref. $^{11}$ we infer $L_{opt} \approx 0.18 \, \mu m$, lower than their PMOKE resolution of 0.28 $\mu m$. We thus conclude that a spatial crossover from the qEW value $\zeta_d \approx 1.25$ to a non-equilibrium exponent $\sim 0.69 \pm 0.07$ must exist. A natural candidate is the Quenched Kardar-Parisi-Zhang (qKPZ) or directed percolation depinning exponent $\zeta_d \approx 0.63$.

![Figure 4](image)

**FIG. 4.** (a) Aspect ratio scaling of $\Delta t = 15$ s WE. The solid (dashed) line shows the expected depinning scaling $S_i \sim L_i^{1+\zeta_d}$ for qEW (qKPZ) class. (b) Scaling of the square width $W^2$ of DW segments of size $L$, for two typical configurations at RT. The solid (dashed) line shows the expected qEW (qKPZ) scaling at depinning, $W^2 \sim L^{2\zeta_d}$, with $\zeta_d = 1.25(0.63)$.

### D. Event lengths

In Fig. 5(a) we show the areas $S_i$ vs the major axis length $L_i$ of each WE. The difference with Fig. 3(a), where $\Delta t = 15\, s$, is that now we plot WE for all the $\Delta t$, from 15s to 180s, in order to observe the effects of large WE. We show both the cloud obtained from raw data and an averaged version by grouping areas in small logarithmic increasing bins and by taking the corresponding average value of $L_i$ in such groups. We compare with the depinning scaling $S_i \sim L_i^{1+\zeta_d}$ expected for cluster events (CE) in the creep regime$^{25}$, both for the qEW class where $\zeta_d = 1.25$ and for the qKPZ class, where $\zeta_d \approx 0.63$. At $L_i \approx L^* = 2 \, \mu m$ a clear crossover is observed (indicated by the vertical line). As can be appreciated in the figure, for $L_i < L^*$ a better agreement is obtained for qEW, as compared for instance with the qKPZ class.

In Fig. 5(b) we show the (non normalized) probability distribution $P_{CE}(L)$ for all the $L_i$ observed. As for $P_{CE}(S)$, we observe a broad distribution. If WE were a part, single, or dominated by single CE we expect indeed WE to display power law distributions similar to the ones observed for depinning avalanches, $P_{CE}(L) \sim S^{-\tau_{CE}}$ and $P_{CE}(L) \sim S^{-\tau_L}$, where $\tau_{CE}$ and $\tau_L$ are related to depinning exponents. The general exponents are well known$^{29}$

$$\tau_{CE} = 2 - (\zeta_d + 1/\nu_d)/(1 + \zeta_d),$$

$$\tau_L = \tau_{CE} (1 + \zeta_d) - \zeta_d$$

where $\zeta_d$ is the depinning roughness exponent and $\nu_d$ the depinning correlation length exponent. For the qEW universality class, we have $\zeta_d \approx 1.25$ and, by virtue of the statistical tilt symmetry$^{30}$, $\nu_d = 1/(2 - \zeta_d) \approx 1.33$. On the other hand $\zeta_d \approx 0.63$ and $\nu_d \approx 1.73$ for the qKPZ class$^{25}$ where the statistical tilt symmetry is broken. This yields $\tau_{CE} \approx 1.11$, $\tau_L \approx 1.25$ for the qEW class, and $\tau_{CE} \approx 1.25$, $\tau_L \approx 1.42$ for the qKPZ class. The effective power law at intermediate $L \lesssim L^*$ (indicated by the vertical line) is roughly consistent with qEW. Unfortunately however, the effective power law observed in Fig. 5(b) is roughly consistent with both classes, unlike Fig. 3(a) which is more consistent with the qEW class.

The crossover at $L^*$, observed in Figs. 3(a), may be associated to the CE coalescence process occurring for large WE$^{31}$. In that case, a WE area can be written as a sum of a given number $N_{CE}$ of CE areas. $S_{CE} \approx \sum_{i=1}^{N_{CE}} S_i$. Since $P_{CE}(S)$ is a broad distribution, the typical WE area is dominated by the largest areas and thus relates to the typical number of CE as $S_{CE} \approx N_{CE}^{1/\tau_{CE}}$. The lateral size of a WE indeed satisfies an inequality $L_{CE} < \sum_{i=1}^{N_{CE}} L_i$, as the fluctuating CE lateral extensions $L_i$ can now overlap. Since $P_{CE}(L)$ is also a broad distribution, we can use the same extreme value argument to estimate $L_{CE} \lesssim N_{CE}^{1/(\tau_L-1)}$. Combining these results we get $S_{CE} \geq L_{CE}^{\nu_{CE}} \lesssim N_{CE}^{1/(\tau_L-1)} \approx L_i^{1+\zeta_d}$. This shows that WE areas should scale with their length approximately as CE in the $L_i < L^*$ regime, as observed in Fig. 3(a). Above $L^*$ however, where large WE become a non negligible fraction of the interface, the last scaling prediction breaks down. In section III C we discuss a simple model that quantitatively accounts for the crossover observed at $L^*$ in the $S_i$ vs $L_i$ plot.
E. Waiting times

The behaviour at large $\Delta t$, where the probability to observe single CE in a WE decreases, is directly related to the behaviour of the large-size $P_{\text{WE}}(S)$ cut-off, $S_{\text{WE}}$, with $\Delta t$. In such regime we can regard each WE area as the sum of a given number $N_{\text{CE}}$ of cluster areas, $S_{\text{WE}} = \sum_{i=1}^{N_{\text{CE}}} S_i$. As $N_{\text{CE}}$ can only grow irreversibly with $\Delta t$, so does the large size cut-off $S_{\text{WE}}$. Naïvely one may think that $S_{\text{WE}}$ should linearly increase with $\Delta t$ because the sum of all WE areas observed in a region of a fixed lateral size $L$ should grow as $L \Delta t$ in a steady-state regime. As shown in Fig. 3(c) we find instead a sub-linear increase $S_{\text{WE}} \sim (\Delta t/t^{*})^{1/2}$. To make sense of this striking observation it is instructive to regard the area $S_{\text{WE}}$ vs. $\Delta t$ as a continuous-time continuous-jump random-walk, with random CE area increments $S_i$ and waiting times $\delta_i$ for the ignition of a new CE, such that $\Delta t = \sum_{i=1}^{N_{\text{CE}}} \delta_i$. If we assume that the $\delta_i$ are distributed according to $\psi(\delta) \sim t^{*\alpha} \delta^{-(1+\alpha)}$, with $0 < \alpha \leq 1$ characterizing the broadness of $\psi(\delta)$, we get $\Delta t \sim t^{*} N_{\text{CE}}^{1/\alpha}$ for the typical number of events $N_{\text{CE}}$ in a $\Delta t$. Since the same heuristic arguments apply for the broadly distributed CE we get $S_{\text{WE}} \sim N_{\text{CE}}^{1/(\tau_{\text{CE}}-1)}$. Combining the two last results we get $S_{\text{WE}} \sim (\Delta t/t^{*})^{\alpha/(\tau_{\text{CE}}-1)}$, which fairly describes our data of Fig. 5(c) if $\alpha/(\tau_{\text{CE}} - 1) \approx 1/2$. Using $\tau_{\text{CE}} \approx 1.11$ we obtain $\alpha \approx 0.05$.

Broad waiting time distributions have been heuristically derived for creep motion\cite{12,13}, borrowing ideas from more general random energy models (see for instance Ref. 13), and also observed numerically close to the depinning threshold\cite{14}. The basic idea is to assume that the EE barrier distribution behaves as $P(U) \sim \exp[-U/U^*]/U^*$ for a large barrier $U$, with $U^*$ a characteristic energy (with $U$ and $U^*$ in units of temperature). If temperature is small enough, the typical time to overcome $U$ is given by the Arrhenius law, $\delta \sim t^{*} \exp[U/T]$, with $t^{*}$ a characteristic time. Changing variables we obtain $\psi(\delta) \sim \delta^{\alpha} \delta^{-(1+\alpha)}$, with $\alpha \sim T/U^*$. Since clustering implies that not all EE have the same $U$ we will argue that the $\delta_i$ corresponds to the special EE that act as CE epicenters. These EE may be associated to the ones allowing to escape from dominant configurations\cite{15}. Two different predictions for $U^*$ and thus for $\alpha$ are found in the literature. In Ref. 12 it is assumed that $U^* \equiv T_d$, with $T_d$ From Table 1 we obtain $\alpha = 0.04$ for $T = 293 K$ and $\alpha = 0.06$ for $T = 323 K$. Both results are in excellent agreement with our data, which gives $\alpha \approx 0.05$. In Ref. 13 on the other hand, the characteristic energy is taken as the optimal nucleous barrier $U^* = T_d(H_d/H)^{\mu}$, with $\mu = 1/4$ for the one dimensional elastic string. The exponent is thus again nonuniversal but now it is also field dependent, $\alpha \approx (T/T_d)(H/H_d)^{\mu}$. From Table 1 we obtain $\alpha \approx 0.02$ both for the two temperatures and their corresponding fields. This value is only slightly below but is again of the order of $\alpha \approx 0.05$ we infer from our measurements. Both predictions are in rough agreement with the empirical $\alpha \approx 0.05$ we obtain from the time-scaling of $S_{\text{WE}}$ we observe in Fig. 3(c) for the two temperatures. It would be interesting to perform a more systematic study as a function of $T$ and $H$ to further test these theories. The previous observations lead us to argue that WE give access not only to the CE area (at small $\Delta t$) but also to the waiting-time statistics (at larger $\Delta t$). As CE start at a seed EE, the $\delta_i$ must be controlled by their energy barrier distribution\cite{35}.

F. Event correlations

In order to further test the connection between WE and CE we have also studied correlations from the spatio-temporal correlations of the registered positions $x_i$ of the $N$ measured WE epicentres. To do that we used the mean square distance $\langle \delta^2 x(T) \rangle \equiv \sum_{i=1}^{N} [x_{i+n} - x_i]^2/N$, which depends only on the temporal separation $T = nt_0$. For non-correlated WE epicentre sequences, $\langle \delta^2 x(T) \rangle$ tends to a constant value $C = (L_0+1)(L_0+2)/2$, where $L_0$ is ap-
proximately the length of the DW in units of the spatial discretization. Figure 6 shows $\langle \delta^2 x(T) \rangle / C$ measured at $T = RT$ and $H = 46.1$ Oe.

FIG. 6. Normalized mean square distance $\langle \delta^2 x(T) \rangle / C$ as a function of $T$, measured at $T = RT$ and $H = 46.1$ Oe.

$T = RT$ and $H = 46.1$ Oe. One can see that even for short $T$ it becomes approximately constant as expected for uncorrelated events (note that $\langle \delta^2 x(T) \rangle / C > 1$ for large $T$ due to an underestimation of the length of DW). We hence conclude that WE are very weakly correlated in sharp contrast with the predicted EE correlations in Ref. 16 and consistent to what is predicted for CE and more generally for depinning avalanches. This observation further confirms our identification of WE with single CE or with coalesced groups of them, for small or large $\Delta t$ respectively.

G. Heuristic model for large WE

Summing up, our results are consistent with the predictions of Ref.16 after identifying the small $\Delta t$ WE with the predicted CE. At large $\Delta t$ WE can not be single CE however, and deviations from the predicted properties for CE are expected. This is already apparent in Fig. 5(a) where large WE display a clearly different length to area aspect ratio than the expected for CE. Moreover, Fig. 5(a) shows a clear crossover from the expected $S_i \sim L_i^{1+\zeta_d}$ CE behaviour to a different behaviour, rather well described by a new power-law, $S_i \sim L_i^{1.4}$. There is no theoretical predictions yet for this crossover so we propose here a simple, heuristic model.

A very simple model can explain the behaviour of $S_i$ vs $L_i$ observed experimentally, shown in Fig. 5(a). The idea is to think WE as the compact objects formed by random deposition of simulated CE, with a lateral sizes $L_i$ sampled from $P_{CE}(L) \sim L^{-\tau_L}$. We can also assume, for simplicity, that the corresponding areas satisfy a deterministic relation $S_i = L_i^{1+\zeta_d}$, assumption that leads automatically to $P_{CE}(S) \sim S^{-\tau_C}$, with $\tau_L$, $\tau_C$ and $\zeta_d$ related by Eqs. 4 and 5. Both assumptions are reasonable approximations according to creep simulation.

To simulate this model we generated such events in the interval $[0, 1]$, sequentially increasing the number of deposited CE. The process starts with one WE which equals the first deposited CE. Adding more CE may produce more WE (specially for small number of deposited CE) or can decrease their number due to the possible coalescence with an existing WE, if the new CE overlaps it. When a coalescence between a new CE and an existing WE takes place, the area of the resulting WE is the sum of the new CE area with the previous area of the WE, but the lenght of the new WE can either remain constant or increase at its left, right or both corners simultaneously.
In more rare cases the new CE can overlap more than one WE. The process finishes when a single WE spans the whole interval, i.e. when the deposited CE percolate the system. To make statistics over many sequences, at this point we reset the simulation and restart adding a first CE into a new actogram. Fig. 7(d) shows actograms corresponding to four runs.

To be concrete, for the simulations we use the values $\zeta_d = 1.25$, $\tau = 1.11$ and $\tau_L = \zeta_d$ corresponding to the 1d qEW depinning class. We sample the epicenter of each CE from a uniform distribution in the interval $[0, 1]$ and its lateral size $l_i$ by $l_i = 0.05\tau(l_{\text{max}}^{\tau+1} - l_{\text{min}}^{\tau+1})^{1/(\tau+1)}$ with $l_{\text{min}} = 10^{-5}$ and $l_{\text{max}} = 1.5$, and $r$ a different uniform random number in the interval $[0, 1]$. This produces a power-law distribution $P_{\text{WE}}(l) \sim l^{-\tau_L}$ with a cut-off at $l_{\text{max}}$. The area of such CE is simply $S_i = l_i^{\tau+\zeta_d}$. Using different parameters yields qualitatively similar results.

We now discuss the results of the model. In Fig. 7(a) we show that the model reproduces the main features observed experimentally in Fig. 5(a). Small WE, below a we show that the model reproduces the main features associated to the lateral acceleration that occurs when the deposited CE percolate the system. To make statistics over many sequences, at the whole interval, i.e. when the deposited CE percolate the [0 1] interval. In Fig. 7(b) we show the probability to percolate as a function of the number of deposited CE. Most of the points in Fig. 7(a) in the $S_i \sim L_i^{1.4}$ regime, belong to states with a high probability to percolate. It is also worth noting that finite size effects, due to the finiteness of the interval and the broad range of the CE lateral size distribution play an important role here. In 7(c) we show that the average anisotropic aspect-ratio $S_i/L_i^{2.25}$ is unity only for a small number of CE, while the aspect ratio increases as large WE tend to completely overlap with most of the new CE and thus increase their areas without modifying its lateral size. Those states are near to percolate but need a rare CE to overlap the voids between the few remaining WE.

The model presented has some unphysical features. In particular, the random deposition process implies, in the long time limit, a growing interface with a non-stationary width. The model describes satisfactorily the crossover in $S_i$ vs $L_i$ observed in the experiments however, so the necessary surface relaxation effects or correlations that may make the width to saturate are not relevant for the regime we aim to describe. In addition, the number of deposited CE does not strictly represent time. Broadly distributed times between random depositions could be easily added however, in order to further test the picture suggested by the experiments. Particularly, to reproduce the area and lateral size distributions as a function of the window time $\Delta t$ found experimentally.

IV. CONCLUSIONS

From our results the following picture emerges. Creep dynamics is driven by EE with a broad size distribution and a large size cut-off controlling the mean velocity. The seed EE that trigger a cascade of extra EE are separated by broadly distributed waiting times. Repeated, this collective process of ignition and correlated growth produce independent CE statistically very similar to depinning avalanches, that may coalesce into larger compact objects. Hence, CE can be truly regarded as “creep avalanches”. The described picture, that drastically changes the naive view of creep motion as independent thermally nucleated displacements, is likely to appear not only in other magnetic films but in the creep regime of other disordered elastic systems in general.

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Appendix A: Samples & Experimental Protocol

Experiments were mainly performed on a Pt/Co/Pt ultrathin ferromagnetic film, a prototypical system which has been the focus of many studies of domain wall motion. The studied sample was a Pt(4.5nm)/Co(0.7nm)/Pt(3.5nm) thin film, with the thickness of each layer indicated in parenthesis. The film was sputter grown at 300 K on etched Si/SiO$_2$ substrate. The magnetic response of this system to an external out-of-plane magnetic field is characterized by a square magnetic cycle with a well defined remnant magnetization, typical of systems with perpendicular magnetic anisotropy.

Polar magneto-optical Kerr effect (PMOKE) microscopy has been used to image magnetic domains after applying magnetic field pulses perpendicular to the film plane. After fully magnetizing the sample in one direction, a short pulse in the opposite direction and of intensity $H = 130$ Oe was first applied in order to nucleate a seed magnetic domain. Then, a second pulse of duration...
t and intensity $H$ was applied to favour the growth of the initial magnetic domain. DW velocity was then computed as the ratio between the linear advance of the DW and the pulse duration, $v = \Delta x / t$. Experiments were performed at $T = RT$ (room temperature) and $T = 50 \degree C$. To measure velocities between $10^{-9}$ m/s and $10^{-1}$ m/s pulses of different amplitude and duration were used. In all the cases, the total number of pulses was 15 or more and the rise time of the pulses was more than one order of magnitude faster than the pulse duration. The shortest pulse was 1 ms and the largest one 1800 s. Due to the spatial resolution of our microscope, for velocities smaller than $10^{-8}$ m/s we observed that there was no difference in the observed DW velocity if the magnetic field was ON or OFF during the image acquisition. In all the cases given the illumination condition, the used shutter time of the camera was 200 ms.

Although we mainly report results for one region of a specific sample, we have made similar measurements in other regions of the same sample and also in a Pt(6nm)/Co(0.2nm)/Ni(0.6nm)/Al(5nm) sample, where the numbers in parenthesis stand for thickness and the ferromagnetic layer consists in a stack of three Co(0.2nm)/Ni(0.6nm) bilayers (for more information about these samples and their domain wall dynamics see Refs. 41 and 42). In both cases the results are consistent with the main universal results reported for the specific region of the Pt/Co/Pt sample in Section III.

FIG. 8. Histograms for the $N_{\text{int}}(t_0, t) = 1151$ WE shown in Fig. 1(c), obtained by comparing consecutive images taken every $\Delta t = t_0 = 15$ s at $T = RT$ and $H = 46.1$ Oe. The histogram with uniform binning is presented in the left panel, while in the right panel a logarithmic binning is used for same WE. Power-law behaviour at small sizes can be described with $P_{\text{we}}(S) \sim S^{-\tau_{\text{WE}}}$ with the distribution exponent $\tau_{\text{WE}} \approx 0.76 \pm 0.1$.

Magnetization reversal events were previously obtained in irradiated Pt/Co/Pt sample identifying between 30 and 50 events depending on field values. In the present work, as we previously anticipated, we were able to obtain a large amount of WE. This represents a quantitative progress in view of the fact that this allows us to perform a deeper statistical description of the data. Figure 8 shows the obtained histogram of the 1151 WE areas shown in Fig. 1(c) by comparing consecutive images taken every $t_0 = 15$ s. Since we are seeking power-law like distributions and their effective exponents, it is convenient to use logarithmic binning. In the right panel of Fig. 8 we use the same WE used in the left panel to build a new histogram. In this case, dividing the number of events per interval by the width of the interval, the probability distribution is obtained. Fig. 8 shows a power-law signature at small size values with a cut-off around 3 $\mu$m$^2$. The distribution is of the form $P_{\text{we}}(S) = S^{-\tau_{\text{WE}}} G_{\text{WE}}(S/S_{\text{cm}})$, where $G_{\text{WE}}$ is the power-law exponent and $S_{\text{cm}}$ the cut-off value such that the function $G_{\text{WE}}(x)$ rapidly decays for $x \gtrsim 1$.

For a proper comparison with theoretical predictions we discarded in the statistical analysis events touching the borders of the region of interest, i.e. the observation region, otherwise their area would be underestimated. This may affect however the tails of the size distribution, corresponding to large events, with a lateral size of the order or larger than the lateral size $L$ of the observation region. For the range of time windows $\Delta t$ we consider WE of lateral size $L$ are extremely rare however. Indeed we observe a $\Delta t$-dependent but clearly $L$-independent cut-off in our distributions, growing as $\Delta t^{1/2}$ (with a temperature dependent prefactor). Such $\Delta t$ dependence is used to estimate the waiting time distribution exponent for cluster “ignition events”. On the other hand, the power-law decay effective exponent of WE, which is also central to our analysis and for the comparison with theory is not sensible to the tails. This justifies our event detection protocol.

To evidence the robustness of our results, in Fig. 9 we show results for the event area distribution measurements done in a different region of the same Pt(4.5nm)/Co(0.7nm)/Pt(3.5nm) sample (upper panel), and for a Pt(6nm)/Co(0.2nm)/Ni(0.6nm)/Al(5nm) sample (lower panel). As can be appreciated, not only the effective power-law decay exponent is similar, but also the time dependence is qualitatively similar to the one of Fig. 3 reported in in Sec. III for the sample and region we have chosen for most of our analysis.

**Appendix B: Domain wall roughness**

In order to have a more direct estimate of the roughness exponent of our DW, and check consistency with our interpretation of the WE statistics, we have computed the single-value displacement field $u_L(x)$ of segments of given sizes $L$ partitioning a larger DW configuration. Here we discuss our practical method. The displacements $u_L(x)$ for each segment are measured with respect to the straight line fitting each segment. This straight line is also used as the $x$-axis to parametrize the displacement field. Such an approach is justified by taking into account that the theoretical description of a directed driven interface assumes that the interface is flat in average in the direction perpendicular to the motion. The field on mag-
length, $W$ expect normal to the DW. Having kinetic DW on the other hand act as a pressure, allways values of $\zeta$ $L$ discretized interfaces of size $N_i(0.6\text{nm})$ text, and obtained for a different Pt(6nm)/Co(0.2nm)/Ni(0.6nm)]$_3$/Al(5nm) sample used to report most our results in the main (upper panel) of the same Pt(4.5nm)/Co(0.7nm)/Pt(3.5nm)

FIG. 9. WE area histograms obtained from a different region (upper panel) of the same Pt(4.5nm)/Co(0.7nm)/Pt(3.5nm) sample used to report most our results in the main text, and obtained for a different Pt(6nm)/Co(0.2nm)/Ni(0.6nm)]$_3$/Al(5nm) sample (lower panel).

We have tested this methodology numerically on large $(\sum x)$ = $x_0$ with $iqx - \zeta$ with a well defined roughness exponent $\zeta$. We do so by superimposing Fourier modes $U(x)$, where $x = 0, 1, ..., L_0 - 1$, with different precise values of $\zeta$. We do so by superimposing Fourier modes $U(x) = \sum_q U_q e^{iqx}$ with $q = 2\pi n/L_0$ ($n = 0, ..., L_0 - 1$), $U_q$ complex hermitian gaussian amplitudes of zero mean, $\langle U_q \rangle = 0$, and variance $\langle |U_q|^2 \rangle \sim 1/q^{1+2\zeta}$ (also known as “self-affine gaussian signals”[19]). This construction assures that the signal $U(x)$ is periodic, $U(x) = U(x + L_0)$, and self-affine with identical spectral and global exponent $\zeta$. Finally, we compute $W^2(L)$ for these interfaces, by the partition procedure previously described. In Fig. 10 we compare $W^2(L)$ vs $L$ with the corresponding scalings for each $\zeta$, averaged over 10 uncorrelated numerically sampled configurations $U(x)$. A good agreement is always obtained if the fit does not include values of $L$ larger than a fixed fraction of the order of the total size $L_0$, so to have a large number of segments and to reduce boundary effects. It is worth noting that our method also allows to accurately measure values $\zeta > 1$, corresponding to super-rough interfaces. This is an advantage over the displacement correlator function.

FIG. 10. Numerical test for the practical implementation used to compute $W^2$ for a DW. We compute $W^2$ for numerically generated self-affine gaussian signals of size $L_0 = 1024$ for several precise values of $\zeta$. We average over 10 samples for each $\zeta$. The solid lines show agreement with the expected $W^2 \sim L^{2\zeta}$. The method allows to measure super-rough cases $\zeta > 1$.

$B(x) \equiv \frac{1}{L_0} \int_0^{L_0} dx_0 \frac{1}{2} \left| u(x + x_0) - u(x_0) \right|^2 /(L_0 - x)$ which gives the correct global $\zeta$, $B(x) \sim x^{2\zeta}$, only if $\zeta < 1$, otherwise it saturates to $\zeta = 1$.

1. S. Lemerle, J. Ferré, C. Chappert, V. Mathet, T. Giamarchi, and P. Le Doussal, Phys. Rev. Lett. 80, 849 (1998).
2. P. J. Metaxas, J. P. Jamet, A. Mougin, M. Cormier, J. Ferré, V. Baltz, B. Rodmacq, B. Dieny, and R. L. Stamps, Phys. Rev. Lett. 99, 217208 (2007).
3. K.-J. Kim, J.-C. Lee, S.-M. Ahn, K.-S. Lee, C.-W. Lee, Y.-J. Cho, S. Seo, K.-H. Shin, S.-B. Choe, and H.-W. Lee, Nature 455, 740 (2009).
4. J. Gorchon, S. Bustingorry, J. Ferré, V. Jeudy, A. B. Kolton, and T. Giamarchi, Phys. Rev. Lett. 113, 077205 (2014).
5. V. Jeudy, A. Mougin, S. Bustingorry, W. Savero Torres, J. Gorchon, A. B. Kolton, A. Lemaître, and J.-P. Jamet, Phys. Rev. Lett. 117, 057201 (2016).
6. R. Díaz Pardo, W. Savero Torres, A. B. Kolton, S. Bustingorry, and V. Jeudy, Phys. Rev. B 95, 184434 (2017).
7. P. Le Doussal, K. J. Wiese, S. Moulinet, and E. Rolley, Europhys. Lett. 87, 56001 (2009).
8. S. Atis, A. K. Dubey, D. Salin, L. Talon, P. Le Doussal, and K. J. Wiese, Phys. Rev. Lett. 114, 234502 (2015).
9. O. Chepizhko, C. Giampietro, E. Mastrapasqua, M. Nourazar, M. Ascagni, M. Sugni, U. Fascio, L. Leggio, C. Malinverno, G. Scita, S. Santucci, M. J. Alava, S. Zapperi, and C. A. M. La Porta, Proc. Nat. Acad. Sci. 113, 5848 (2016).
11408 (2016).

10 S. S. P. Parkin, M. Hayashi, and L. Thomas, Science 320, 190 (2008).

11 T. Nattermann, Phys. Rev. Lett. 64, 2454 (1990).

12 L. B. Ioffe and V. M. Vinokur, J. Phys. C 20, 6149 (1987).

13 G. Durin, F. Bohn, M. A. Corrêa, R. L. Sommer, P. Le Doussal, and K. J. Wiese, Phys. Rev. Lett. 117, 087201 (2016).

14 L. Laurson, X. Illa, S. Santucci, K. Tore Tallakstad, K. J. Måloj, and M. J. Alava, Nat. Commun. 4, 2927 (2013).

15 V. Repain, M. Bauer, J.-P. Jamet, J. Ferré, C. Chappert, and H. Bernas, Europhys. Lett. 68, 460 (2004).

16 E. E. Ferrero, L. Foini, T. Giamarchi, A. B. Kolton, and A. Rosso, Phys. Rev. Lett. 118, 147208 (2017).

17 P. Chauve, T. Giamarchi, and P. Le Doussal, Phys. Rev. B 62, 6241 (2000).

18 See Supp. Mat. in Ref. 16.

19 J.-P. Bouchaud and A. Georges, Physics Reports 195, 127 (1990).

20 V. H. Purrello, J. L. Iguain, A. B. Kolton, and E. A. Jagla, Phys. Rev. E 96, 022112 (2017).

21 We show both the original noisy $S_i$ vs. $L_i$ curve and a smoothed curve, obtained by grouping similar areas in small bins and by assigning the average of $L_i$ to each area bin. Indistinguishable results are obtained by binning $L_i$ rather than $S_i$.

22 E. E. Ferrero, S. Bustingorry, and A. B. Kolton, Phys. Rev. E 87, 032122 (2013).

23 A. Rosso and W. Krauth, Phys. Rev. Lett. 87, 187002 (2001).

24 L.-H. Tang, M. Kardar, and D. Dhar, Phys. Rev. Lett. 74, 920 (1995).

25 A.-L. Barabási and H. E. Stanley, Fractal Concepts in Surface Growth, Cambridge University Press ed. (Cambridge, 1995).

26 The downward deviation crossover for large WEs or lateral acceleration is due to coalescence effects in the proximity of a percolation transition. See Sec [H] for further details.

27 A. B. Kolton, A. Rosso, T. Giamarchi, and W. Krauth, Phys. Rev. Lett. 97, 057001 (2006).

28 A. B. Kolton, A. Rosso, T. Giamarchi, and W. Krauth, Phys. Rev. B 79, 184207 (2009).

29 S. Zapperi, P. Cizeau, G. Durin, and H. E. Stanley, Phys. Rev. B 58, 6353 (1998).

30 M. Kardar, Phys. Rep. 301, 85 (1998).

31 In our protocol, the maximum WE lateral size $L_i$ was limited by the lateral size $L_{roi}$ of the region of interest. We did not consider WE touching nor spanning completely the region of interest. If spanning WE were considered we would have $S_{WE} \approx v \Delta t L_{roi}$ in the long-time limit, with $v$ the average velocity, shown in Figs. 1(a),(b).

32 V. M. Vinokur, M. C. Marchetti, and L.-W. Chen, Phys. Rev. Lett. 77, 1845 (1996).

33 C. Monthus and T. Garel, Phys. Rev. E 78, 041133 (2008).

34 E. A. Jagla, unpublished.

35 Barrier distributions for dominant metastable states have been computed in Ref. [28] for forces below and near $f_c$. For small systems a roughly exponential right tail followed by a cut-off can be appreciated.

36 L. Foini, A. Rosso, private communication.

37 S. Emori and G. S. D. Beach, J. Phys. Condens. Matter 24, 024214 (2012).

38 K.-W. Moon, D.-H. Kim, S.-C. Yoo, C.-G. Cho, S. Hwang, B. Kahng, B.-C. Min, K.-H. Shin, and S.-B. Choe, Phys. Rev. Lett. 110, 107203 (2013).

39 A. Hrabec, N. A. Porter, A. Wells, M. J. Benitez, G. Burnett, S. McVitie, D. McGrouther, T. A. Moore, and C. H. Marrows, Phys. Rev. B 90, 020402 (2014).

40 A. W. J. Wells, P. M. Shepley, C. H. Marrows, and T. A. Moore, Phys. Rev. B 95, 054428 (2017).

41 J.-C. Rojas-Sánchez, P. Laczkowski, J. Sampaio, S. Collin, K. Bouzehouane, N. Reyren, H. Jaffrès, A. Mougin, and J.-M. Georges, App. Phys. Lett. 108, 082406 (2016).

42 N. B. Caballero, I. Fernández Aguirre, L. J. Albornoz, A. B. Kolton, J. C. Rojas-Sánchez, S. Collin, J. M. George, R. Diaz Pardo, V. Jedy, S. Bustingorry, and J. Curiale, Phys. Rev. B 96, 224222 (2017).

43 A. Rosso, R. Santachiara, and W. Krauth, Journal of Statistical Mechanics: Theory and Experiment 2005, L08001 (2005).