Barkhausen avalanches in anisotropic ferromagnets with 180° domain walls

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We show that Barkhausen noise in two-dimensional disordered ferromagnets with extended domain walls is characterized by the avalanche size exponent $\tau_s = 1.54$ at low disorder. With increasing disorder the characteristic domain size is reduced relative to the system size due to nucleation of new domains and a dynamic phase transition occurs to the scaling behavior with $\tau_s = 1.30$. The exponents decrease at finite driving rate. The results agree with recently observed behavior in amorphous Metglas and Fe-Co-B ribbons when the applied anisotropic stress is varied.

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Barkhausen noise measured in disordered ferromagnets at low temperatures under the condition of slow driving by an external field through the hysteresis loop exhibits scaling behavior without tuning of any parameter [1–7]. The scaling properties of Barkhausen avalanches obtained in various alloys can be grouped in three distinct universality classes differentiated by the value of the avalanche size exponent as [8] $\tau_s \approx 1.3$, $\tau_s \approx 1.5$, and $\tau_s \approx 1.7$. However, the origin of the scaling of Barkhausen noise (BN) and the occurrence of different universality classes is still not fully understood.

The above-mentioned measurements [1–7] are done mostly on thin ribbon samples (thickness $d < 200$ μm). The variety of measured scaling exponents can be attributed to the differences in the applied driving conditions and to the diversity of domain structures occurring in different samples. It was found that the exponents decrease continuously with increasing driving rate [2,5,6]. The structure of domains in commercial alloys has been studied by a variety of techniques (see, for instance, [9]). In thin ribbons of the amorphous Metglas Fe-B-Si [1,7,9] and Fe-Co-B alloys [5,6], which are annealed in a parallel field [9] or under an applied anisotropic stress [6,7], a structure with few domains occurs with 180° domain walls parallel to the anisotropy axis. The demagnetizing fields, which depend on the form of the sample, and the range of interactions play an important role for the equilibrium domain structure as well as for the domain-wall dynamics [3,5]. It has been demonstrated that the domain structure in the amorphous Metglass [7] and Fe-Co-B alloys [6] can be controlled by varying tensile stress and that short-range interactions dominate over dipolar forces [6]. In addition, in these systems the demagnetizing fields are minimized with longitudinal anisotropy [9]. Hence, in our study we can neglect dipolar forces, but we take into account the existence of domain walls. Numerical studies of BN using short-range Ising models with various types of disorder [10–12] (see also [13] for a more realistic type of interactions) usually start from a uniform ground state and nucleate a random pattern of clusters of reversed spins by increasing the external field, thus neglecting the preexisting domain structure.

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In this paper we simulate Barkhausen avalanches using a model with preexisting *extended* domain walls which are confined in two dimensions, motivated by the stress-induced anisotropy in realistic systems [6,7]. Increasing the external magnetic field the domain wall may either move through a random medium or new domains may nucleate when this is energetically favorable [14]. It has been suggested recently that interface depinning in a random medium is responsible for the scaling behavior of BN [3,5,7]; however, models of an elastic interface yield scaling exponents that are lower than the measured ones [3,7]. In this work we use a ferromagnetic model with short-range interactions and a random-field pinning and we show that the results compare well with two universality classes of measured scaling exponents. In addition, a dynamic phase transition between these two scaling behaviors appears when the strength of disorder (or size of the domains) is tuned. It has been recognized recently that disorder effects are enhanced leading to smaller domain sizes by decreasing either tensile stress [6,7] or grain sizes [13]. Motivated by these suggestions, here we study the influence of disorder on the scaling properties of Barkhausen noise. We expect that the results are relevant to realistic samples in which the domain size exceeds the sample thickness.

We consider an Ising model on a square lattice of size $L \times L$ assuming that local random fields h_i are generated by coarse-graining from an original disorder [14] in the presence of the external magnetic field H:

$$\mathcal{H} = -\sum_{\langle i,j \rangle} J_{i,j} S_i S_j - \sum_i (h_i + H) S_i . \qquad (1)$$

Here we set $J_{i,j} = 1$ to be a constant interaction between nearest-neighbor spins $S_i = \pm 1$. Hence, all fields and energies are measured in units of $J_{i,j}$. A Gaussian distribution of h_i is assumed with zero mean and width f. We create an initial domain wall in the $\langle 11 \rangle$ direction by a rotation of the lattice by $\pi/4$ and setting all spins except those in the first row opposite to the external field [15]. Periodic boundaries in the direction of the interface and fixed boundaries in the perpendicular direction are applied. The dynamics consists of a spin alignment parallel to the external field when the local field $h_i^{loc} \equiv \sum_j J_{i,j} S_j + H + h_i$ exceeds zero. For the simulation of a hysteresis loop the field updates are adjusted

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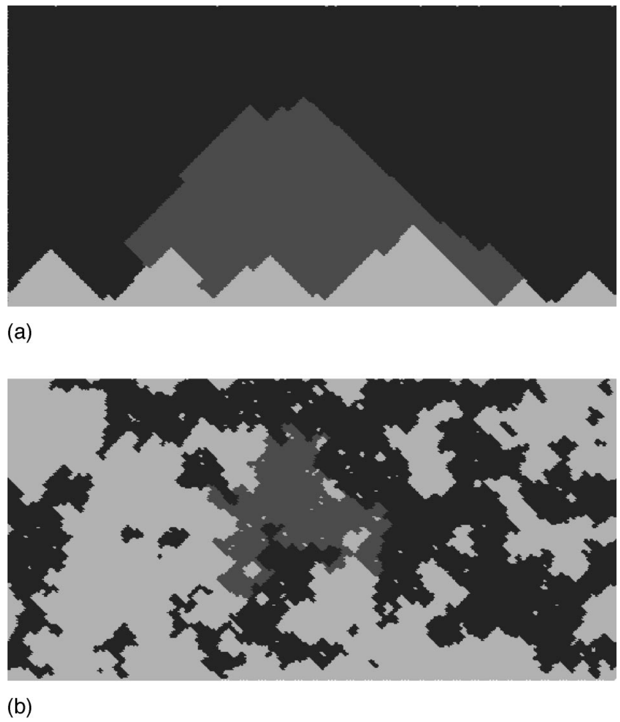

FIG. 1. Emerging cluster structure of flipped (bright) and unflipped (dark) spins, and the most recent cluster (gray), illustrating self-affine growth at low disorder (a), and nucleation and blocking by previous clusters at high disorder (b). $L=200$.

to the minimum local field (infinitely slow driving), thus the driving field is uniform in space but fluctuates in time. We also briefly discuss the effects of finite driving rates.

It should be stressed that at each time step we update *all spins*, which is suitable for a globally driven magnetic system. In this way new domains are nucleated when it is energetically allowed, in contrast to models of driven interfaces, where an update is restricted only to the sites located next to the interface, resulting in ''percolation'' growth even at high disorder [16]. Another important feature of our model is the anisotropy [17] due to the initial conditions: the threshold driving forces in the parallel and perpendicular directions appear to be different. This model has the following advantages: (1) In the limit of vanishing disorder the $\langle 11\rangle$ interface can be moved by an infinitesimally small field. This bypasses the problem of threshold energy $2J$, which occurs in a $\langle 10\rangle$ interface [16] implying that a large lattice size have to be used in order to find a spin with a random field large enough to surmount the energy barrier. Therefore, here we can apply smaller lattice sizes and vary disorder configurations (we use up to 2×10^3 configurations and up to $L=400$). (2) The $\langle 11\rangle$ interface depins at infinitesimally small field at low disorder $f\rightarrow 0$. Hence, an *upper limit of disorder* $f^\star$ *exists* at which depinning is no longer possible, and nucleation of new domains in the interior becomes favorable at large enough fields. This feature comprises an important difference compared to the model of Ref. [10], where nucleation of a single spanning cluster at $f\rightarrow 0$ requires energy threshold $4J$, and thus becomes obscured by finite lattice size and large fields.

Figure 1 shows snapshots of simulated systems for high and low disorder. For low disorder only domain wall motion occurs. For values of f that exceed a certain critical value $f^\star$

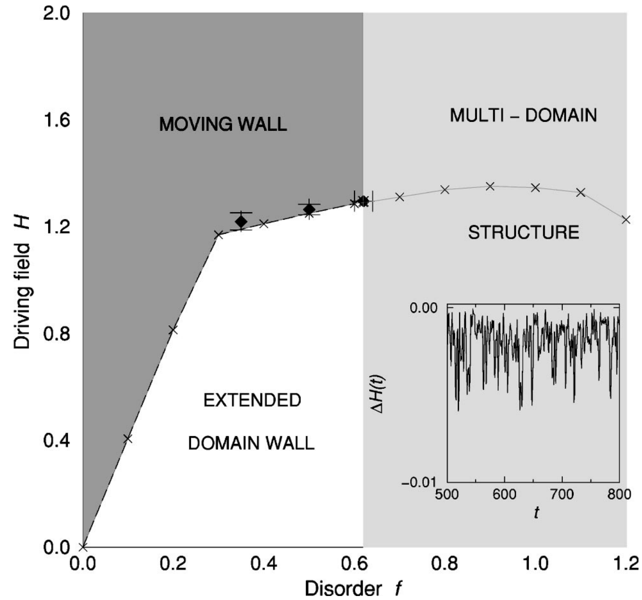

FIG. 2. Phase diagram in the plain disorder–driving field (both in units of J). The vertical dashed line at $f^\star=0.62$ separates the high and low disorder region. Long-dashed line with crosses: coercive fields $H_0(f)$; filled symbols: critical fields $H_c(f)$ obtained from scaling fits. Inset: Time [in Monte Carlo steps (MCS)] series of driving field increments ΔH (in J/MCS) for $f=1.2$.

(determined below), domains nucleate inside the system, thus leading to the same structure of clusters as in systems without an initial domain wall [12,19]. These two regions are shown in the phase diagram in Fig. 2 together with simulation results for the coercive fields $H_0(f)$ of the hysteresis and the critical fields $H_c(f)$ of the depinning transition (explained below).

We apply the quasistatic driving described above (an example of time series of field increments is given in the inset to Fig. 2) and monitor the motion of domain walls. The number of flipped spins s between two consecutive locally stable configurations of the wall determines the size of Barkhausen avalanche. The avalanche size distribution $D(s,f)$ is shown in Fig. 3 for various values of disorder f. Taking the avalanche statistics along the ascending part of the hysteresis loop (until eventually depinning occurs) and averaging over many disorder configurations we find the slopes according to $D(s)\sim s^{-\tau_s}$ as $\tau_s=1.54$ for $f\leqslant 0.6$, and $\tau_s=1.30$ above $f^\star\approx 0.6$. The cutoff also decreases with f. A similar behavior was found experimentally in Metglas 2605TCA in Ref. [7]. With the stress σ varying in the range from 0–525 MPa the slopes of the size distributions are reported to vary from 1.29 to 1.60, and the cutoffs increase [7]. In $Fe_{64}Co_{21}B_{15}$ the measured slope was 1.28 [6] for σ up to 140 MPa. The applied anisotropic stress stretches the domain walls, thus for the degree of disorder f we have $f\sim 1/\sigma^x$. We also measure the distributions of the linear extensions of avalanches in the directions parallel (w) and perpendicular (h) to the wall leading to the anisotropy exponent $\zeta\equiv(\tau_w-1)/(\tau_h-1)=0.92$. (See Table I.)

In the low disorder region the domain wall depins when the driving field exceeds a critical value $H_c(f)$. To investigate the depinning transition we start with a flat wall and apply a constant field H measuring the *long-time limit* of the

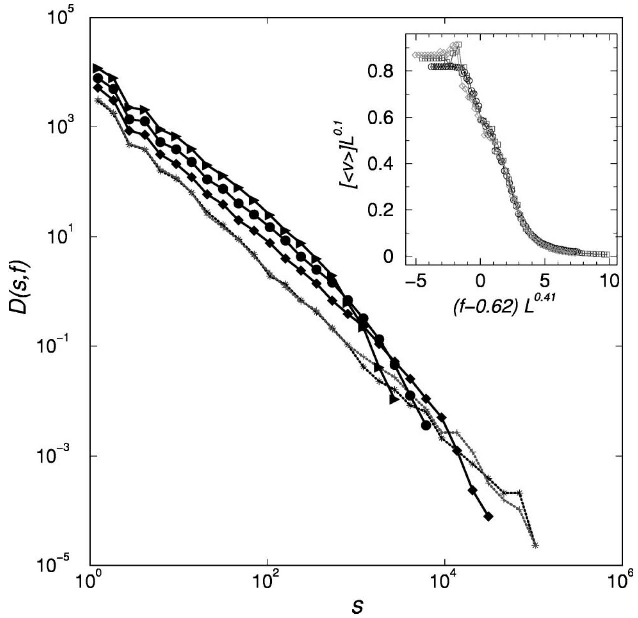

FIG. 3. Probability distribution of avalanche size $D(s,f)$ vs size s for various values of disorder f=0.4, 0.6, 0.9, 1.0, and 1.1 (bottom to top). Inset: Scaling collapse according to Eq. (2) of the velocity of domain wall $[\langle v\rangle]$ vs f for L=128, 196, and 256. Each point is averaged over 200 samples.

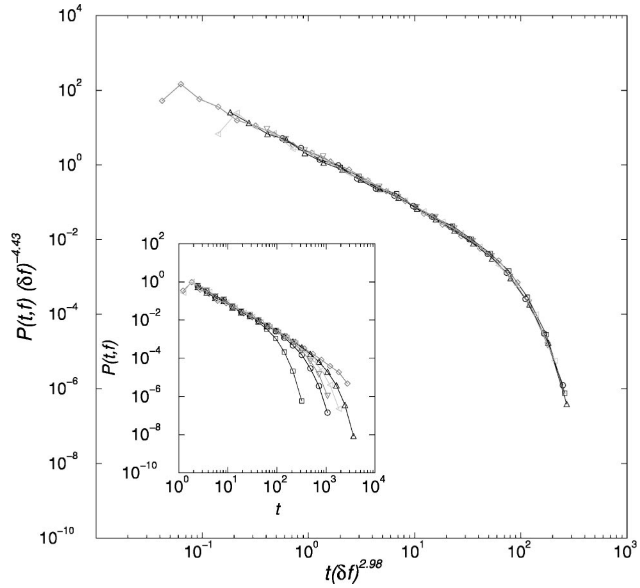

FIG. 4. Inset: Probability distribution of durations of Barkhausen avalanches $P(t,f)$ vs duration t in MCS, for $L=400$ and for $f=$ 1.2, 1.0, 0.96, 0.92, 0.88, and 0.84 (left to right). Main figure: Scaling collapse of the data (see text).

domain wall velocity v, which is defined as the number of flipped spins *per internal time step* relative to L. The velocity averaged over 2000 disorder realizations, $[v]$, is the order parameter of the depinning transition. Here $[v]$ and its fluctuations are obtained for $L=100$, 200, and 400 and the critical exponents and critical fields $H_c(f)$ (also shown in Fig. 2) are determined by a finite size scaling plot of the general form

$$Y(X,L)=L^{-Z_Y}\mathcal{Y}(L^{1/\nu}X) \qquad (2)$$

with $X\equiv[H-H_c(f)]/H_c(f)$ (see Table I, left side, and Ref. [19]). Here $Z_Y\equiv\beta/\nu$ for the analysis of the order parameter $[v]$, and $Z_Y\equiv\gamma/\nu$ for the analysis of the fluctuations of $[v]$. Our results are compatible with the correlation length exponent $\nu=1.23\pm0.04$, $\beta=0.43\pm0.03$, and $\gamma=1.52\pm0.06$. The exponents are universal in the region $0.35\leqslant f\leqslant 0.6$. Using the scaling relation $\beta/\nu=z-\zeta$ valid for the depinning transition [17] we find the dynamic exponent $z=1.27$. Then the relation $D(\tau_s-1)=z(\tau_t-1)$, where the fractal dimension of anisotropic avalanches is $D=1+\zeta$, leads to the duration exponent $\tau_t=1.83$. These results are in good agreement with universal criticality in a class of driven dynamical systems recently discussed in Ref. [20]. The fraction of active sites at time t scales as $N(t)\sim t^{\kappa}$, with $\kappa=D/z-1=0.51$, compared with 0.58 in [21]. The value for the exponent ν compares well with one in [16,18]. For the (local) roughness exponent ζ values in the literature ranging from 0.5 to 1.23 can be attributed to the influence of the anisotropy [17], elasticity of the interface [22], distribution of disorder [15,16], and driving conditions [20,23]. Further analysis, e.g., by the dynamic renormalization group, is necessary in order to determine if the value $\zeta=0.92$ found in this work represents a new universal behavior or a crossover due to finite size effects. This value suggests closeness to the class of models with vanishing interface velocity [18] and self-organized depinning [23]. At low disorder lateral motion of the interface makes it possible to overcome strong pinning centers and to maintain the self-affine growth (see Fig. 1). However, at the transition we find $\phi\equiv2\nu(1-\zeta)-\beta=-0.23$, thus $\lambda_{eff}\to0$ resulting in a different critical behavior compared to the case $\lambda_{eff}\to\infty$ studied in Refs. [17,18].

TABLE I. Scaling exponents obtained by *direct* simulation of various quantities for the low (LD) and high (HD) disorder region. Error bars estimated from independent fits are within ±0.03.

Exponent	Value (LD)	Exponent	Value (HD)
τ_s	1.54	τ_s	1.30
τ_w	2.03	τ_t	1.47
τ_h	2.12	D	1.88
ζ	0.92	z	1.23
β/ν	0.35	β/ν	0.06
γ/ν	1.22	γ/ν	2.26
$1/\nu$	0.82	$1/\nu$	0.42

In order to determine the largest disorder where depinning is still possible, $f^{\star}$, we can formally extend the above study of the domain wall velocity now averaged over hysteresis loop and disorder $[\langle v\rangle]$, and the corresponding susceptibility. Applying then Eq. (2) with $X\equiv(f-f^{\star})/f^{\star}$ (scaling collapse is shown in the inset to Fig. 3) we find $f^{\star}=0.61\pm0.02$ and the exponents β, γ, and ν shown in Table I, right side.

In the region of high disorder $f>f^{\star}$ nucleation of new domains of finite size which are blocking each other's spatial extent becomes the dominant feature that determines the scaling properties of BN [see Table I (right)]. The section of power-law behavior of the distribution of avalanches increases with decreasing f, in qualitative agreement with increasing stress in experiments [6,7]. In the inset to Fig. 4 the distribution of avalanche durations is shown for various values of f above $f^\star$. The main Fig. 4 shows the scaling plot $P(t,f)=(\delta f)^{z\nu\tau_t}\mathcal{P}(t(\delta f)^{z\nu})$, where $\delta f \equiv (f-f^\star)/f^\star$, obtained by using $f^\star=0.62$ and the products of the exponents $z\nu\tau_t=4.43$ and $z\nu=2.98$. From a similar scaling collapse of the size distribution we find $D\nu\tau_s=5.61$ and $D\nu=4.30$. Thus, these values together with the ones in Table I (right side) lead to $\nu=2.3\pm0.1$, $\beta=0.12\pm0.03$, and $\gamma=5.2\pm0.1$. Note that the value of $f^\star=0.62$ within statistical error bars ± 0.03, which we estimated from two different types of scaling fits (see Figs. 3 and 4) is by no means definitive. (The error bars are expected to increase when more scaled quantities or wider range of system sizes are explored.) However, since no exact results are available, this value can be regarded as a rough estimate of the critical disorder.

By applying a finite driving rate $r \equiv \Delta H/H_{max}$, where H_{max} is the saturation field, the cutoffs of the distributions in the high disorder region increase (see also [12]), whereas slopes of the distribution decrease with r due to mainly the coalescence of avalanches. For the size of avalanches we find, for instance for $f=0.88$, $\tau_s=1.18$ for $r=0.01$ and $\tau_s=1.10$ for $r=0.02$.

In conclusion, we have shown that the anisotropic two-dimensional motion of domain walls pinned by quenched impurities and short-range interactions are relevant for the scaling behavior of Barkhausen noise in the presence of extended domain walls. We find two universality classes with $\tau_s=1.54$ and $\tau_s=1.30$, in a good agreement with experiments in amorphous Fe-B-Si and Fe-Co-B ribbons under anisotropic stress. By slow driving the scaling behavior is robust in a wide range of the driving field and disorder (or stress) values. These two universality classes correspond to the motion of extended domain walls (i.e., at low disorder or high stress), and many finite domains (high disorder or low stress), respectively. The scaling exponents decrease with finite driving rate. We find that the domain structure changes via a dynamic phase transition at a critical disorder, which can be directly monitored in experiments by tuning uniaxial stress.

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