Domain Size Effects in Barkhausen Noise

M. Bahiana, Belita Koiller
Instituto de Física, UFRJ
Rio de Janeiro, RJ, Brazil

S.L. A. de Queiroz
Instituto de Física, UFF
Niterói, RJ, Brazil

J. C. Denardin, R. L. Sommer
Departamento de Física, UFSM
Santa Maria, RS, Brazil

The possible existence of self-organized criticality in Barkhausen noise is investigated theoretically through a single interface model, and experimentally from measurements in amorphous magnetostrictive ribbon Metglas 2605TCA under stress. Contrary to previous interpretations in the literature, both simulation and experiment indicate that the presence of a cutoff in the avalanche size distribution may be attributed to finite size effects.

The Barkhausen effect consists of magnetic noise caused by erratic jumps in the magnetization of a ferromagnetic material, under an increasing applied magnetic field [1]. A simple explanation for this effect is the combination of random pinning of domain walls by defects and the driving external field, which is essentially the same mechanism present in stick-slip processes [2]. Recently the statistical behavior of Barkhausen noise has attracted much interest, due to the possibility of providing an experimental realization of self-organized critical (SOC) behavior [3–5]. The subject is controversial, however. We concentrate here on the results obtained by Urbach et al. (UMM) [4] and Perković et al (PDS) [5].

UMM measured the avalanche size probability distribution function in an Fe-Ni-Co alloy, and found power-law decay over approximately two decades, followed by an exponential cutoff. The same result was also observed in numerical simulations of the interface motion. The power-law behavior, obtained without any intentional fine-tuning of parameters, suggests that this system self-organizes into a critical state. On the other hand, PDS argued that such behavior can be explained without resource to SOC concepts: in their view, the power-law decay followed by a cutoff is evidence that the system is near but not quite at a conventional critical point. They performed simulations for the random-field Ising model (RFIM) under an external field, taking the local (pinning) fields to be gaussian-disordered with standard deviation $R$. The avalanche-size distribution is also characterized by a power law followed by a cutoff, and the power-law regime increases over several decades as $R$ approaches a critical disorder $R_C$.

Although UMM and PDS approach the problem with apparently similar models, their conclusions regarding the critical nature of the Barkhausen noise are in contradiction. Here we show that in reality, the ingredients used in either model differ in crucial aspects where the onset of SOC is concerned, so it is not surprising that they end up with different findings.

We investigate this question by using the simple model proposed by UMM [4] for the motion of a single domain wall in the Barkhausen noise regime. We find that the existence of a cutoff in the UMM model can be traced back to finite-size effects; experimental results, also to be described, bear out the idea that the cutoff to be found there originates from corresponding aspects in real systems.

In UMM’s model, the interface at time $t$ is described, in space dimensionality $d$, by its height $h(\tilde{\rho}_i,t)$, where $\tilde{\rho}_i$ is the position-vector of site $i$ in a $(d-1)$-dimensional lattice. At each $t$, the height function $h_i = h(\tilde{\rho}_i,t)$ is assumed to be single-valued, so there are no overhangs on the interface. Thus the interface element corresponding to the $d$-dimensional position-vector $\tilde{r}_i = (\tilde{\rho}_i,h_i)$ may be unambiguously labelled by $i$. Simulations are performed on a $L^{d-1} \times \infty$ geometry, with the interface motion set along the infinite direction. Therefore finite-size effects are controlled by the length parameter $L$. Each element $i$ of the interface experiences a force of the form:

$$f_i = u(\tilde{r}_i) + \frac{k_z}{2} \left[ \sum_{j=1}^{z} h_{ij} - z h_i \right] + H_e , \quad (1)$$

where
The magnetization is updated, and this process continues until
we concentrate on the 3-d results. After a transient, the effective field always settles onto a critical value which depends on \( R \), \( k \) and \( \eta \). The fact that the system tunes itself to a constant effective field is in itself an indication of SOC-like behavior. If the external field is started at a higher value, a large avalanche occurs and brings the effective field back to the adequate value. On the other hand, if the field starts from zero there is a transient corresponding to a series of small avalanches. We find that the number of avalanches in this transient is proportional to \( L^{d-1} \). This means that an infinite system would need an infinite number of avalanches to reach criticality. As an illustration of this behavior, we present in Figure (1) the effective field in a 3−d system for \( R = 5.0, k = 1 \) and \( \eta = 0.05 \) for different simulation cell sizes \( L \).
FIG. 1. Simulation on a $3 - d$ lattice for systems with $R = 5.0$, $k = 1$, $\eta = 0.05$ and different widths. Starting the external field at $H = 0$, the effective field $H_e$ grows linearly and saturates at a critical value after a transient. Here we see this behavior as a function of the avalanche number and, in the inset, the same but with avalanche number scaled by $L^2$.

The calculated avalanche-size distribution corresponds to a power-law with a cutoff, as obtained by Urbach et al in Barkhausen experiments and in simulations and also by Perković et al within their RFIM model. It was shown in the latter reference that the cutoff is intrinsic to the RFIM, if the system is away from the critical point $R_C$. In what follows we present evidence that the nature of the cutoff in the Barkhausen effect, both in simulations of the Urbach model and in experiments, is a finite-size effect. The characterization of the cutoff as a finite-size effect was recently suggested by Narayan through the analysis of a continuum model closely based on UMM’s, though no attempt was made there to quantify the relationship. We have examined the dependence of the cutoff on the simulation cell size $L$ by collecting series of 100,000 avalanches for $L = 50, 80, 150, 200$ and 400. Figure 2 shows the avalanche size distribution for some values of $L$ in $3 - d$ lattices. The transient was eliminated by starting the external field near $H_c$. We can clearly see that the cutoff increases with $L$. The histograms were fitted by the function $P(A) \propto A^{-\alpha} \exp(-A/A_0)$ with $\alpha$ in the range $(1.23 \pm 0.02 - 1.35 \pm 0.02)$. The parameter $A_0$, which defines the cutoff size, is strongly dependent on $L$: $A_0 \propto L^{1.4\pm0.1}$. We have also varied the disorder parameter $R$, in an attempt to find an effect similar to that described by PDS in their RFIM model. However, the final picture was qualitatively always the same as shown in Figure 2, with roughly the same power-law dependence of $A_0$ on $L$. In $d = 2$ we used lattices of width $L = 100, 500, 1000, 2000, 3000$ and 5000 from which a value of $\alpha = 0.83 \pm 0.08 - 1.03 \pm 0.03$ and $A_0 \propto L^{0.78\pm0.06}$ were obtained. Thus, we conclude that the presence of a cutoff in the avalanche size histogram is a finite-size effect. Though it may seem puzzling that the finite transverse dimensions can influence the characteristics of interface motion along the unbound direction of growth, a qualitative picture of the corresponding mechanism is as follows. Each time the $(d - 1)$-dimensional lattice is swept, the number of distinct chances for the interface to move is of order $N \sim L^{d-1}$; thus, if the typical probability for a given interface element not to advance (that is, to have $f_i < 0$) is $p$, the interface as a whole will come to a halt, marking the end of an avalanche, only if $f_i < 0$ for all
elements. This happens with probability $p^N$, hence for large $L$ larger avalanches become less unlikely.

The above result clearly demonstrates the different nature of the cutoff in the UMM and PDS simulation models. Both models describe domain walls advancing in a disordered magnetic medium: The key difference between them is the presence of a demagnetizing field, proportional to the growing domain magnetization. This field is an essential element to bring the simulations into an SOC regime.

Although variations in $L$ are easily accessible in simulations, this is not a trivial parameter to vary in experiments. For magnetic systems, one would expect $L$ to be related to the typical magnetic domain size in a sample. Magnetic materials have the interesting property of magnetostriction, which is the change in internal domain configurations as a response to applied anisotropic stress. Positive magnetostrictive materials show an increase in internal domain wall lengths under applied stress [9], and may therefore be employed in investigating experimentally $L$-dependent effects in Barkhausen noise experiments. We performed measurements in a disordered material with this property, namely amorphous magnetostrictive ribbon Metglas 2605TCA under stress $\sigma$. Different stress values were applied in order to increase domain wall length, as seen in ref. [9]. Stress is expected also to change the domain wall thickness $\delta = \sqrt{A/K}$, where $A$ is the effective exchange constant, $K = \frac{3}{2}\lambda_S \cdot \sigma$ and $\lambda_S = 27 \cdot 10^{-6}$ is the saturation magnetostriction constant for this material. In amorphous materials, fluctuations in the stress due to local composition fluctuations generate the pinning sites for domain walls. Their magnitudes are much larger than those originated from the external forces. Thus, the effect of the applied stresses will be to order the domain wall structure by changing the domain size and length. The magnetoelastic anisotropy, on the other hand, tends to stretch existing domain walls [9].

Metglas 2605TCA samples (80 mm $\times$ 1 mm $\times$ 30 $\mu$m) were pre-annealed in an Ar flow at 300°C for 15 min. in order to decrease the stress level associated to the fabrication process. The measurements were performed in an open magnetic circuit. As a consequence, there is a global effective field acting on the domain walls; also, the average magnetization rate $\dot{M}(t)$ and differential susceptibility were kept constant. The samples were cycled in their hysteresis loops, excited by a slowly varying (triangular, 0.2 Hz) field. The Barkhausen signal was detected by a small (approx. 5 mm) coil wound around the central part of the sample. The signal was preamplified by a SR550 low pass amplifier and digitized by a TDA320 oscilloscope. The low pass amplifier was set to an upper frequency limit equal to half the sampling frequency. The waveform generator, current source and preamp were fed by batteries in order to increase the signal to noise ratio. At each cycle and starting from a given value of a trigger field, a time series was acquired and stored for further processing. For each stress level, we identify an avalanche with a jump in the voltage level $V$. As in Ref. [4], a threshold voltage was established according to experimental resolution. This threshold defines the low-end cutoff of the avalanche distribution, and shall not concern us particularly here as it is the high-end of the distribution that will be sensitive to finite-size effects. It is to be noted, however, that especially for low-stress data, the combination of experimental resolution available and the intrinsic characteristics of Metglas resulted in a rather narrow power-law region. This in turn yielded quite a large spread in the fitted values of the effective power $\alpha$, as seen below.

The avalanche (voltage jump) size distributions were obtained for $\sigma = 0, 17, 30, 80, 100, 150, 180, 230, 300, 400$.
and 525 MPa. Some of these results are shown in Fig. 3.

![Histogram of voltage jumps for different values of applied stress $\sigma$.](image)

The dashed line, included for comparison with Fig.2 and with UMM’s experimental data, has the same slope as the dashed line there, $-1.3$.

It is clear that the cutoff increases with the applied stress. Since stress increases the magnetic domains, what we see here is again the dependence of the cutoff on a typical domain size, in accordance with the simulations results. The experimental histograms where also fitted with the function $P(V) \propto V^{-\alpha} \exp(-V/V_0)$ giving $\alpha$ in the range $0.29 \pm 0.1 - 1.6 \pm 0.1$ and $V_0 \propto \sigma^{0.65 \pm 0.04}$. The fact that experimental values of $\alpha$ fall in a similar range to those from the 3–d simulations is both to be expected and in line with the earlier results of UMM. On the other hand, the power that governs the dependence of $A_0$ on $\sigma$ need not coincide with that which relates $A_0$ to $L$ in the numerical work. In order to predict a relationship between the two quantities, one would need to work out the connection of the physical mechanisms driving the interplay between finite transverse dimensions and avalanche sizes, in the interface model, to the corresponding ones between applied stress and domain wall length in actual samples. Thus far, we have not been able to do so.

From the above results we conclude that the cutoff in the avalanche size histogram in Barkhausen systems is a finite size effect. This, together with the presence of a self-tuning effective field and negative time correlation for short times, is a strong indication that SOC is present. Also, we find that any attempt to model Barkhausen noise must necessarily include the demagnetizing field, for this is the key ingredient for the above mentioned self-tuning.

**ACKNOWLEDGMENTS**

This work was partially supported by CNPq, CAPES and FAPERGS(Brazil). We thank M. Novak, J. Urbach and A. Hansen for interesting discussions.

[1] H. Barkhausen, Phys. Zeits 20, 401 (1919).
[2] H. J. S. Feder and J. Feder, Phys. Rev. Lett. 66, 2669 (1991).
[3] K. P. O’Brien and M. Weissman, Phys. Rev. E 50, 3446 (1994).
[4] J. Urbach, R. Madison, and J. Markert, Phys. Rev. Lett. 75, 276 (1995).
[5] O. Perković, K. Dahmen, and J. Sethna, Phys. Rev. Lett. 75, 4528 (1995).
[6] H. Ji, M.O. Robbins, Phys. Rev. A 44, 2538 (1991).
[7] P. Cizeau, S. Zapperi, G. Durin, and H. Stanley, Phys. Rev. Lett. 79, 4672 (1997).
[8] O. Narayan, Phys. Rev. Lett. 77, 3855 (1996).
[9] K. Yamamoto, T. Sasaki, and Y. Yamashiro, J.Appl.Phys. 81, 5796 (1997).