SHAPE OF A BARKHAUSEN PULSE

F. Colaiori\textsuperscript{a} S. Zapperi\textsuperscript{a} G. Durin\textsuperscript{b}

\textsuperscript{a}INFN-SMC, Dipartimento di Fisica, Universit\`a "La Sapienza", P.le A. Moro 2 00185 Roma, Italy
\textsuperscript{b}Istituto Elettrotecnico Nazionale Galileo Ferraris, strada delle Cacce 91, I-10135 Torino, Italy

Abstract

The average shape of the pulse in Barkhausen noise has been recently proposed as a tool to compare models and experiments. We compute theoretically the pulse shape of Barkhausen noise in a model describing the motion of a domain wall in an effective Brownian potential. In this framework, the pulse shape is related to the properties of the excursion of a random process in a $\log(x) - kx$ potential. We record the Barkhausen noise in polycrystalline FeSi materials, and compare the pulse shape with the one predicted by the domain wall model.

Key words: random magnets, hysteresis, Barkhausen noise, random walk

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The Barkhausen noise has been incessantly investigated because of its practical application and theoretical implications. Experiments show that both size and duration of avalanches of spin reversal are power law distributed over several decades. Recently, the average pulse shape has been proposed as a sharper tool to test models against experiments [1].

The scaling analysis suggested in Ref. [1] is based on a simple relation between the average size $\langle S \rangle$ of an avalanche and its duration $T$ scaling as $\langle S \rangle \sim T^{1/\sigma_{\nu z}}$ where $1/\sigma_{\nu z}$ is a combination of critical exponents, defined in Ref. [1]. The average avalanche has an universal shape given by $v(t,T) \sim T^{1/\sigma_{\nu z} - 1} g(t/T)$, where $g$ is a universal scaling function. Similarly, when considering the signal $v$ as a function of the magnetization $s = \int_0^t vdt$, one gets $v(s,S) \sim S^{1-1/\sigma_{\nu z}} f(s/S)$ where $f$ is another scaling function. In the model presented in Ref. [1], $g$ is found to be an inverted parabola, but this result does not fit with experiments.

The Barkhausen noise can be described in terms of a phenomenological model, known as ABBM [2], which describes the wall as a rigid interface in an effective Brownian pinning potential. The crucial assumption of a Brownian correlated pinning potential was done on phenomenological basis, since this kind of correlation is observed in experiments, but the model can be obtained as a mean-field version of a more general flexible domain wall model [3]. Considering the domain wall velocity as a function of magnetization $s$, one obtains the following Langevin equation [3]

$$\frac{dv}{ds} = \frac{c}{v} - k + \eta(s)$$

(1)

where $\eta$ is uncorrelated Gaussian noise with $\langle \eta(s)\eta(s') \rangle = 2\delta(s - s')$, and $c$ is a dimensionless parameter proportional to the applied field rate. If we neglect the contribution of the demagnetizing factor $k$, Eq. 1 describes the motion of a 1d random walk in a logarithmic potential $E(v) = -c \log(v)$, where the magnetization $s$ plays the role of time.

The corresponding Fokker–Plank equation is

$$\frac{\partial P(v,s)}{\partial s} = \frac{\partial}{\partial v} \left( \frac{-c}{v} + \frac{\partial}{\partial v} P(v,s) \right),$$

(2)

where $P(v,s)$ is the probability to find the walk in $v$ at $s$. We are interested in a solution of this equation with the initial condition $P(v,0) = \delta(v - v_0)$ and an absorbing boundary at the origin $P(v = 0,t) = 0$. This solution can be expressed in terms of modified Bessel functions [5]. For the interesting case $0 < c < 1$, which is the condition of the
ABBM model to have power laws in the avalanche distribution [2], the probability $P(v, s|v_0, 0; c)$ for a walk starting at $v_0$ to be at $v$ after a “time” $s$, in the limit $v_0 \to 0$ is simply proportional to a power of $v$ times a Gaussian with variance $s$, and the average excursion is the ratio between two successive moments of a Gaussian with variance $w = 2s(S-s)$, and it is thus simply proportional to $\sqrt{w}$. The normalized average excursion is therefore given by $\langle v \rangle \propto \sqrt{s(S-s)}$. It is also possible to calculate the function $g$. By definition the avalanche size $s$ at time $t$ is given by the integral of $v(t, T)$ from time zero to time $t$:

$$s = \int_0^t dt' v(t', T) \propto T^{1/\sigma \nu z} \int_0^{t/T} g(x) dx,$$

which provides an expression of $s = s(t, T)$ as a function of $t$ and $T$. Imposing $v(t, T) = v(s = s(t, T), S(T))$ gives an integral equation for $g$ involving $f$: $g(x) \propto f \left( \int_0^x g(x') dx' \right)$. Using the form of $f$, we can solve this equation with the boundary conditions $g(0) = g(1) = 1$. The solution is $g(x) \propto \sin(\pi x)$. Summarizing, for the normalized avalanche we obtain

$$v(s, S) = S^{1-\sigma \nu z} \pi \sqrt{(s/S)(1-s/S)},$$

$$v(t, T) = T^{1/\sigma \nu z-1} \pi /2 \sin(\pi t/T).$$

It is interesting to compare the theoretical results given above with the experimental average shapes. In Figs. 1 we plot the signal voltages both as a function of time $v(t, T)$ and magnetization $v(s, S)$ rescaled using the theoretical value, $1/\sigma \nu z = 2$. The demagnetizing fields correspond to a bias to the random walk, which has the effect of introducing a cutoff in the distribution of avalanche sizes and durations, but has no effect on the shape of the scaling functions. Despite the asymmetry and the deviation from scaling observed in the experimental shapes, the predictions based on the ABBM model are in reasonable agreement with experiments, supporting the validity of the interface model approach.

References

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