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Single-wall dynamics and power law in bistable magnetic microwires

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Abstract
The dynamics of a single domain wall in finite glass-coated microwires is reported in the low-field regime below the switching field. The power law of the single domain wall propagating over a large distance is confirmed. Three regions are determined for the propagating domain wall. Below some critical field, H_0, the domain wall is pinned at the wire's end. Just above that critical field, the wall moves in the adiabatic regime, interacting with the defects during its propagation, with an average velocity of $v = S'(H - H_0)^\beta$. At high field, the domain wall propagates in the viscous regime and its average velocity is proportional to the applied magnetic field, H. An analysis of the temperature dependence of the scaling factor β is further reported. It is also shown that the domain wall mobility parameter S' is field-independent and is proportional to the domain wall mobility S in the viscous regime.

1. Introduction

Magnetic domain walls in small magnetic devices are used to encode information for storage, for sensing in some magnetic sensors or to perform logic operations [1]. A full understanding of the magnetization reversal process in small structured systems is the key for future applications in the above-mentioned devices (hybrid integrated circuits, race-track memory and sensors) [2, 3]. Although, the domain wall dynamics has been extensively explored theoretically, there is a lack of experimental work, with some provided only in the last few years [1, 3] and some of it sometimes controversial [5, 6]. Understanding and controlling the domain wall propagation through the real material containing the defects can improve both the dynamic behaviour of the wall through these structures and its application functionality.

We present the study of the domain wall dynamics on glass-coated microwires in the low-field regime, where a universal power law is observed. Magnetic glass-coated microwires prepared by the Taylor–Ulitovski method exhibit unique domain structure resulting from shape and magnetoelastic anisotropy [4–6]. Particularly, that of positive magnetostriction Fe based microwires consists of a large single domain with magnetization oriented axially, surrounded by a radially oriented domain structure. Moreover, small closure structures appear at the ends of the wire to reduce the stray field energy. As a result, the axial magnetization process runs through the depinning and propagation of a single domain wall from one closure domain at one end. Although having almost bulk structure, glass-coated microwires show a unique nearly ideal magnetization reversal process, characterized by its fast velocity or high mobility [7, 8], the investigation of which reveals much of interest to help comprehend what is occurring in smaller nanostructures. Particularly, a very fast domain wall has been observed in microwires [9] that can reach even supersonic velocities 20 000 m s^{-1} [10]. One of the parameters that are responsible for the fast domain wall propagation is a negative critical propagation field [9] that would lead to the virtual propagation in a negative field [8]. In order to reach a deeper understanding of the reversal process, a new experiment related to domain wall dynamics in the low-field regime has been carried out.

The domain wall dynamics is governed by a linear dependence of the domain wall velocity v on the applied field H [11]:

$$v = S(H - H_0), \quad (1)$$

where S is the domain wall mobility and H_0 is the critical propagation field, which should be overcome to observe

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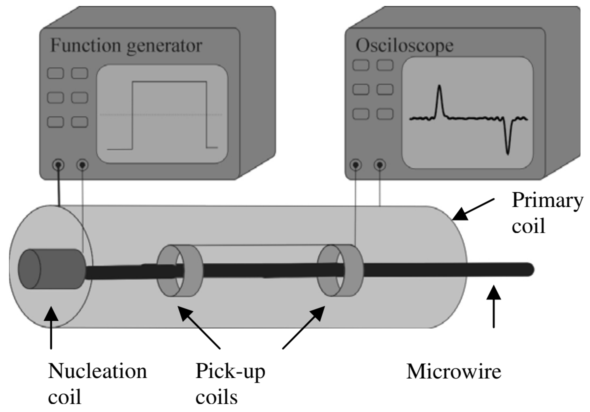

Figure 1. Schematic diagram showing the experimental setup for the domain wall dynamics measurements.

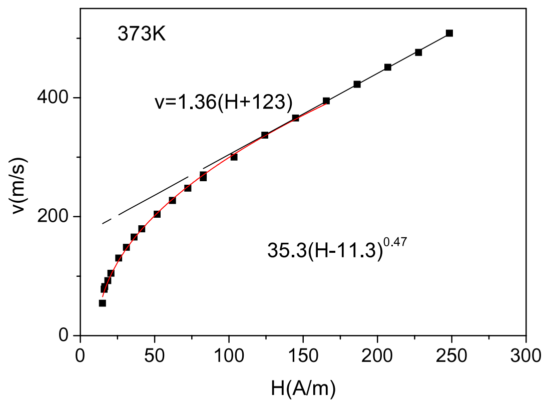

Figure 2. Domain wall velocity v as a function of magnetic field amplitude H.

(This figure is in colour only in the electronic version)

the domain wall propagation. Such linear dependence has been experimentally confirmed many times in different materials [12–14]. However, recent measurements of the single domain wall propagation in glass-coated microwires [7, 15] show unusual dynamics where extrapolation of linear behaviour leads to negative critical propagation field, H_0. To obtain further information on the process, measurements of the low-field domain wall propagation dynamics have been performed.

2. Experimental details

A classical Sixtus and Tonks [16] like experiment is used to determine the domain wall dynamics on a $Fe_{77.5}Si_{7.5}B_{15}$ amorphous microwire, 11 μm in diameter, coated by a 9 μm thick Pyrex layer. The system consists of four coaxial coils (figure 1). The primary coil, 10 cm long, generating the exciting field was fed by 30 Hz frequency AC square current creating a homogeneous field along the wire, 10.5 cm long, that can be taken as static during wall propagation. A small nucleation coil located at the very end of the wire was used to generate an additional local magnetic field of the same frequency, in phase with the field in the primary coil and higher than the switching field, that is, high enough to overcome the pinning field of the closure domain. Two secondary coils, symmetrically disposed at the centre of the primary coils and separated by $L = 6$ cm, are connected in series opposition so, two sharp opposite peaks are picked up at an oscilloscope upon passing the propagating wall. The coils system allows us to identify the propagating wall direction the velocity of which is calculated as $v = L/t$, where t is the time between two maxima in the *emf* recorded peaks. The system is placed inside a specially designed cryostat system, the details of which can be found elsewhere [7, 8], enabling the measurement in the temperature range from 77 to 380 K.

The local field generated by the nucleation coil (1 cm long) fully reverses magnetization within a small region at one end of the wire, which quickly enlarges the action of the homogeneous field created by the solenoid. The amplitude of the nucleation field was kept just above the critical propagation field H_0 (given by equation1) small enough not to disturb the experiment, but high enough to ensure the initialization of propagation.

3. Results and discussion

In previous works, the expected linear behaviour according to equation (2) has been measured under homogeneous fields stronger than the so-called switching field. Here, we also pay particular attention to the case of propagation for applied field below the switching field.

Below the static switching field of the closure domain wall, a characteristic deviation from the linearity is found (see figure 2). Alternatively, the low-field domain wall dynamics can be described by a power law:

$$v = S'(H - H_0')^{\beta} \tag{2}$$

where S' is an effective domain wall mobility parameter, H_0' is the dynamic coercive field and β is the power exponent. Such a power law results from the interaction of the propagating wall with the defects of the material, different sources of which in the actual amorphous microwire have been described before [8, 17]. As was shown in [7] and [17], the domain wall potential E consists of two terms: the long-range magnetoelastic one and the short-range terms arising from the pinning of the domain wall on the defects in the amorphous medium. The pinning centres are randomly distributed along the amorphous microwire and the pinning field. Therefore, the domain wall potential fluctuates as the domain wall moves through the material, and the restoring force α acting on the domain wall due to the gradient of the internal potential is given by $\alpha = \mathrm{d}E/\mathrm{d}x$. When the domain wall passes the region with a local maximum of the restoring force α, a local jump occurs until the wall reaches a new site with the restoring force α greater than the force $2M_S H$ acting on the domain wall.

Therefore, the domain wall motion in the low-field limit is adiabatic. Under the action of a small force $2M_S H$, the domain wall moves slowly close to some local minimum in its potential that arises from the elastic interactions within the domain wall and from the impurities pinning. At some

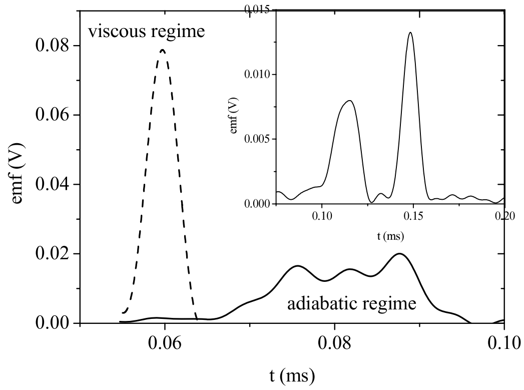

Figure 3. The emf recorded waveform measured in the viscous and adiabatic regime (at very low domain wall velocity) at 273 K. The inset shows how the peak is separated into two peaks when the pinning is so strong that the domain wall stops inside the pick-up coil (below H_0').

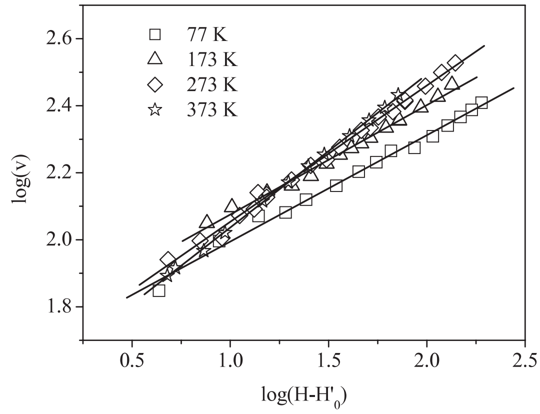

Figure 4. Domain wall velocity as a function of magnetic field amplitude for a range of indicated measuring temperatures. Full lines represent the linear fit.

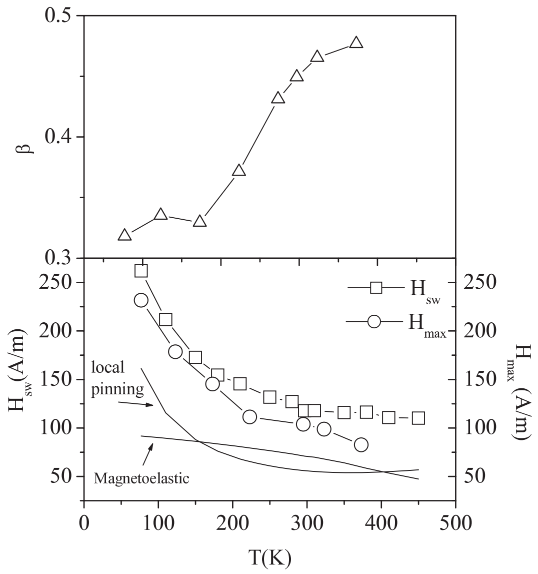

Figure 5. Temperature dependence of the scaling exponent β (up). Temperature dependence of the switching field H_{sw} and its contributions taken from [5] and theoretically estimated H_{max} according to equation (6) (down).

point, the local minimum disappears and the domain wall moves forward rapidly to another local minimum. Such motion is characterized by intermittent jumps from defect to defect. Hence, the local domain wall dynamics description is governed by the generalized equation (1):

$$v = S(H - (H_{dm} + H_p)), \tag{3}$$

where H_{dm} includes all long-range contributions (such as the geometry dependent demagnetizing field and the magnetoelastic contribution) and is separated from a random component, H_p, that includes all short-range counterfield contributions. The pinning field, H_p, is assumed to exhibit statistical properties governed by details of the local pinning potentials that inhibit domain wall motion. In the first approximation, the distribution of the pinning field takes Gaussian shape with the width of R. Then the magnetization change during the domain wall jump is given by the power law [18]:

$$\Delta M \sim ((R - R_c)/R_c)^{\beta} \tag{4}$$

where R_c is a critical distribution width below which the small intermittent domain wall jumps do not appear. As a result, the domain wall moves with average velocity ($v = \Delta M/\Delta t$) given by equation (2) [18–21].

The different shape of the induced peaks at the pick-up coils (obtained by a 'single shot' acquisition) also confirms the presence of the two regimes (figure 3). The perfect symmetric shape measured in the viscous regime confirms the planar shape of the domain wall, which propagates at constant velocity. The emf waveform measured in the adiabatic regime (at very low domain wall velocity -50 m s^{-1}) suggests the fluctuation of the domain wall velocity during its motion across the randomly distributed defects. When the applied field is very close to the critical field H_0', the domain wall can even stop in the centre of the pick-up coil (see inset of figure 3). Then one can observe two peaks that correspond to two Barkhausen jumps, similarly to the experiments in [22]. The domain wall can depin from the defect due to the thermal activation [23, 24]. In this case it does not propagate along the entire wire (6 cm) since it finally remains pinned somewhere in the middle.

The observed nonlinear dynamics at low fields in figure 2 is not surprising. It was first calculated by Nakatani *et al* for thick sub-micron strip line [25], and assumed theoretically from the measurement of Yang and Erskine [26]. Here, we confirm the scaling behaviour of the single domain wall during its propagation over large distances, which has been taken as a manifestation of the nonlinear dynamics and criticality in complex systems [27].

The power law is universal and valid on a wide range of scales where crackling noise is detected [28–33]. The power law at low fields is here confirmed also in a wide temperature range, as proved in the log–log plot in figure 4. However, the power exponent is temperature dependent (see figure 5) and increases from $\beta = 0.32$ at 77 K up to $\beta = 0.48$ at 373 K. This evolution can be treated in terms of the change of the domain wall shape. We can assume two contributions to the domain wall potential in amorphous microwires [7, 17], firstly, the long-range magnetoelastic one that changes with temperature due to the different thermal expansion coefficient of metallic nucleus and glass coating [34, 35]. A short-range contribution comes from the interaction of the domain wall with the defects on the atomic level. Due to the interaction energy with local spontaneous magnetization these mobile defects try to align into the most favourable orientation. Amorphous alloys have a low packing density because of the steric misfit between atoms of different atomic radii. Therefore, even at low temperature small rearrangements of atoms are possible by jumps of atoms into the neighbouring free volumes. The total interaction energy of the domain wall with the mobile defects, also called stabilization energy, can be expressed as [36]:

$$E_p = \frac{2}{15}\frac{\rho_0}{kT}\varepsilon_p^2\left[-2\delta_0 + \frac{x^2}{\delta_0}\right]F(t,T), \qquad (5)$$

where ρ_0 is the density of the mobile defects, ε_p corresponds to the interaction energy of the mobile defects with the local spontaneous magnetization, k is a Boltzmann constant, T is the temperature, δ_0 is the domain wall width, x is the domain wall position and $F(T,t)$ is a relaxation function: $F(T,t) = (1 - e^{(-t/\tau)})$, where t is the time of measurement and τ is the relaxation time, given by the Arrhenius equation: $\tau = \tau_0 e^{Q/kT}$, τ_0 being a pre-exponential factor and Q denoting the activation energy of the mobile defects. It was shown earlier [7, 8, 17] that pinning of the domain wall on the mobile defects becomes an important mechanism especially at low temperatures. When the temperature decreases, the mobile defects lose their mobility, increase the local anisotropy and the pinning of the domain wall on such defects is stronger.

Hence, one possible explanation of the temperature dependence of β arises from the domain wall shape change due to pinning. At high temperatures, the long-range magnetoelastic contribution prevails, the domain wall pinning on the local defects is small (see figure 5—contributions to the switching field) and the domain wall prefers to keep a planar (rigid) shape in order to decrease the stray fields. According to the random field theory [19–21], the coefficient $\beta \to 1/2$. At low temperature, the pinning forces from the randomly distributed defects are much stronger than elastic forces and the interface breaks up [20]. The domain wall motion at any point depends on its pinning on the local defects and the domain wall takes a flexible shape. This results in the decrease of the power exponent β down to 0.32 at 77 K. This value corresponds well with the random field model [20, 21]. These results are also consistent with the temperature dependence of the switching field and its contributions (see also figure 5) [7]. Moreover, such a temperature dependence of β does not occur in microwires, where the pinning contribution to the hysteresis mechanism is negligible [37]. Unfortunately, there is no direct experimental evidence for the domain wall shape change. The propagating domain wall in microwires is shielded by the radial domain structure so the direct observation of the domain wall is impossible.

Table 1. Fitted parameters of the domain wall dynamics to equations (1) and (2).

T (K)	S'	H_0' (A m^{-1})	β	H_t (A m^{-1})	S (m^2 A^{-1} s^{-1})	H_0 (A m^{-1})
77	47.9	17.1	0.318	216	0.51	−317
123	50.7	12.4	0.335	154	0.66	−251
173	56.3	10.5	0.329	124	0.84	−197
223	50.4	7.8	0.371	82	0.97	−176
273	40.4	5.4	0.431	103	1.14	−151
295	35.1	13.0	0.449	110	1.07	−147
323	35.2	12.4	0.465	113	1.15	−150
373	35.3	11.3	0.477	124	1.36	−123

Another important parameter is the effective domain wall mobility parameter S' given in equation (2). Its role is not clear and there is also speculation that the nonlinear domain wall dynamics can be treated in terms of the field dependence of the mobility parameter S', keeping the power exponent $\beta = 1$ [19]. Anyway, according to equations (2) and (4), $v \sim ((H - H_0')/H_0')^\beta$. Then, the mobility parameter S' is equal to $S^*/(H_0')^\beta$. The new parameter S^* is proportional to the domain wall mobility in the viscous regime S given by equation (1), where the proportionality constant, H_{max}, has the dimensions of magnetic field ($S^* = S.H_{max}$). As observed in figure 5, H_{max} takes similar values to the switching field obtained in [5]. In addition, it has the same temperature dependence, which points to the fact that both fields are governed by the same mechanism. Finally, we have found the relation between the mobility parameter S' given by equation (2) and domain wall mobility in the viscous regime S as:

$$S' = S.H_{max}/H_0^\beta, \qquad (6)$$

which supports the fact that the domain wall mobility can be taken as a constant and it is not field dependent. At least in the adiabatic and viscous regime.

At higher fields, larger than the fluctuations of the wall potential, equation (2) transforms into equation (1). In this viscous regime, above H_t, when the wall propagates at constant velocity it is not locally pinned so it propagates in a single continuous step without interaction with local defects. Accordingly, equation (2) can be modified as:

$$v = v(H_t) + S(H - H_t). \qquad (7)$$

An analytical description in the whole range is:

$$v = \begin{cases} 0 & \text{for } H < H_0' \\ (S.H_{max}/H_0'^\beta)(H - H_0')^\beta & \text{for } H_0' < H < H_t \\ v(H_t) + S(H - H_t) & \text{for } H_t < H \end{cases} \qquad (8)$$

where the fitted parameters are collected in table 1.

4. Conclusions

In conclusion, the experimental evidence of the theoretically predicted power law dependence of the domain wall velocity on low applied magnetic field is reported here for the case where the domain wall moves slowly, interacting with the defects presented in the materials. We offer one possible explanation of the temperature dependence of the power exponent β in terms of the domain wall shape change. The power exponent β approaches values of $1/2$ in the high temperature regime, indicating the rigid domain wall. At low temperatures, the domain wall pinning on the defect prevails and the domain wall becomes flexible, resulting in the decrease of β. We also show that the mobility parameter is proportional to the domain wall mobility for the viscous regime and to the field $H_{\max}$ at which the domain wall dynamics changes from adiabatic to viscous. This also shows that the nonlinear dynamics cannot be treated in terms of the field dependence of the domain wall mobility.

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