Switching-field distribution in amorphous magnetic bistable microwires

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We present the temperature dependence of the switching field distribution of amorphous magnetic bistable microwires, which has unusual behavior in comparison with the experimental data presented up to now. The distribution is solved in terms of the thermoactivated model and its unusual temperature dependence is explained by the two contributions to the domain wall pinning in the amorphous microwires. The distribution width is found to be proportional to the switching field.

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I. INTRODUCTION

An analysis of the micromagnetism of the magnetization reversal process of mesoscopic magnetic materials with a simple domain structure is of basic interest as it allows one to gain information about its intrinsic magnetic and thermodynamics characteristics. While magnetization rotation processes are mainly determined by the strength of magnetic anisotropies considered at long range scale, in the case of wall displacements, short range scale ordering/disordering characteristics may play a decisive role. Thus, nucleation, propagation, and depinning phenomena are connected to short range ordering relaxation, disaccommodation effects. Most commonly, a large number of walls are involved in the reversal process, and only in a few cases a relatively simple reversal process model can be modeled, as, for example, in the classical case of iron whiskers[1] as a consequence of the well defined domain structure determined by magnetocrystalline anisotropy.

Alternative nearly ideal materials are considered in the present study: positive-magnetostriction amorphous microwires.[2] Their quite simple domain structure, determined by magnetostrictive and shape energy terms, consists of a single-domain inner core with magnetization oriented parallel to the wire axis that is surrounded by the outer domain shell with the magnetization oriented perpendicular to the wire axis.[3] Only small closure domain structures appear at the end of the microwires in order to decrease the otherwise very large stray fields energy.[4] This domain structure results in the specific magnetization process when the wire is magnetized in a field parallel to its axis. Then, at the switching field, the magnetization reversal process runs in a single giant Barkhausen jump as a domain wall depins from the closure structure at one of the ends and propagates along the entire microwire.[5,6]

In contrast with previous attempts to study the reversal process of other relatively simple domain structures,[7,8] the outstanding case of bistable amorphous microwires gives us the possibility to study the magnetization process by a single domain wall displacement without any interaction with other walls. Due to the intrinsic fluctuations, the measured switching field shows distributed values instead of being single valued, as has been quite recently reported.[9] This distribution can give us valuable information about intrinsic thermodynamics of the domain wall movement in amorphous materials.

In this contribution, a detailed model is developed for the description of the switching field distribution characteristics in amorphous alloys microwires as well as for its temperature dependence.

II. THERMOACTIVATED MECHANISM MODEL

As mentioned above, the magnetization process of magnetically bistable amorphous microwires runs by a single giant Barkhausen jump between two stable remanent states. Therefore, the depinning of a domain wall from the closure structure at one end is responsible for the coercivity mechanism.

Before the action of the external magnetic field, the position x of the pinned domain wall is given by its potential $W(x)$ minimum which in amorphous materials is mainly determined by the magnetoelastic potential.[10]

Under the action of the applied magnetic field, H (see Fig. 1), the total free energy $G(x)$ of the domain wall is given by[11]

$$G(x) = W(x) - 2\mu_0 M_s HAx, \tag{1}$$

where μ_0 is permeability of vacuum, M_s is saturation magnetization and A is the area of the domain wall assumed to be rigid. When the applied field, H_3, reaches the value of the switching field, H_{sw}, the energy barrier exists no more and

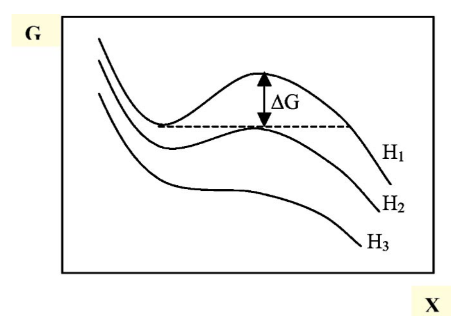

FIG. 1. Dependence of the free energy G of the closure domain wall on its position x at different external magnetic field $H_1 < H_2 < H_3$.

the domain wall jumps into the more favorable position. But if the external field H_1 is lower than the switching field, the domain wall lays in a minimum of $G(x)$ until it can overcome the maximum $G(x_{max})$ either by thermal fluctuations (when at high enough temperatures) or by tunnelling (at low temperatures).[12]

In order to obtain the dependence of the energy barrier ΔG on the external magnetic field H, let us call x_0 the equilibrium position of the domain wall in the zero external magnetic field. The potential of the domain wall can be expanded as a function of the position x close to $x_0(\Delta x = x - x_0)$ as[13]

$$W(x) = W_0 + W_0' . \Delta x + \frac{1}{2}W_0''(\Delta x)^2 + \frac{1}{6}W_0'''(\Delta x)^3, \quad (2)$$

where due to Δx smallness, 4th and higher power terms will be neglected. At the zero applied magnetic field, $H=0$, the domain wall stays at position $x=x_0$ where the potential wall is minimum:

$$W_0' = 0, \quad (3)$$

The giant Barkhausen jump occurs when the switching field, H_{sw}, is reached at

$$\left.\frac{\partial G}{\partial x}\right|_{H=H_{sw}} = 0 \quad (4)$$

and

$$\left.\frac{\partial^2 G}{\partial x^2}\right|_{H=H_{sw}} = 0. \quad (5)$$

From Eqs. (3)–(5) it follows that

$$W'' = \frac{k^2}{\chi_0} \quad (6)$$

and

$$W''' = -\frac{k^3}{2\chi_0^2 . H_{sw}}, \quad (7)$$

where χ_0 is the initial susceptibility of the domain wall at the zero external field and $k = 2\mu_0 M_s A$.

The dependence of the energy barrier on the external magnetic field can be calculated by inserting Eqs. (2), (3), (6), and (7) into Eq. (1). In the local extreme of the free energy,

$$\left.\frac{\partial G}{\partial x}\right|_{H<H_{sw}} = 0. \quad (8)$$

The energy barrier given by the difference between the local maximum and minimum of the free energy is

$$\Delta G(H) = G_{max}(H) - G_{min}(H), \quad (9)$$

and taking into account Eqs. (1), (8), and (9), it can finally be expressed as

$$\Delta G(H) = \frac{8}{3}\chi_0\sqrt{H_{sw}}(\Delta H)^{3/2}, \quad (10)$$

where $\Delta H = H_{sw} - H$.

The 3/2-power law has been also found when calculating the energy barrier dependence on the external magnetic field.[14,15]

Due to thermal fluctuations of the switching field, its probability, dp, lies within the range between ΔH and $\Delta H + d(\Delta H)$ and is given by[16]

$$dp = A \exp(-\Delta G/kT)d(\Delta H), \quad (11)$$

where A is a constant, and kT is a thermal energy.

Inserting the field dependence of the energy barrier Eq. (10) into Eq. (11) and introducing the reduced fluctuation $\Delta h = \Delta H / H_{sw}$, a linear dependence of the logarithm of the probability density $\ln(dp/d(\Delta h))$ on the $(\Delta h)^{3/2}$ is found:

$$\ln(dp/d(\Delta h)) = \beta - \alpha(\Delta h)^{3/2}, \quad (12)$$

where

$$\alpha = (8/3)\chi_0 H_{sw}^2/kT, \quad (13)$$

$$\beta = (2/3)\ln(\alpha) - \ln(2) - \ln[y(\alpha^{1/3})], \quad (14)$$

and

$$y(\alpha^{1/3}) = \int_0^{\alpha^{1/3}} z \exp(\chi^{1/3})\mathrm{d}z, \quad (15)$$

As later analyzed, the linear dependence in Eq. (12) will be very useful for the evaluation of the switching field distribution. Parameter α gives us information about the shape of the domain wall potential. If the switching field value H_{sw} is measured simultaneously with the distribution, the initial susceptibility χ_0 can be obtained which gives us information about the shape of the domain wall potential bottom.

III. SWITCHING FIELD DISTRIBUTION

The switching field distribution was measured on the amorphous glass-coated microwire of the composition $Co_{68}Mn_7Si_{15}B_{10}$ prepared by the Taylor-Ulitovsky method. The microwire was 17 mm long, with the diameter of metallic nucleus of 8 μm and a total diameter of 20 μm.

The switching field distribution was obtained by evaluating 2000 consecutive switching processes. The measurements were performed within the temperature range between 77 and 350 K. Further details on the experimental technique can be found elsewhere.[9]

The switching field distribution measured at different temperature is shown in Fig. 2. As observed, the distribution becomes narrower and the switching field takes smaller values as the measuring temperature increases. This is in opposition to previous works[13,15,17] where a 2/3-power dependence of the switching field distribution on the temperature was found. In order to solve this discrepancy, a two-potential model is considered for the domain wall responsible for coercivity mechanism of bistable amorphous microwires.

IV. DISCUSSION

A. Temperature dependence of mean switching field

As recently introduced,[9] the temperature dependence of the switching field distribution can be interpreted in terms of

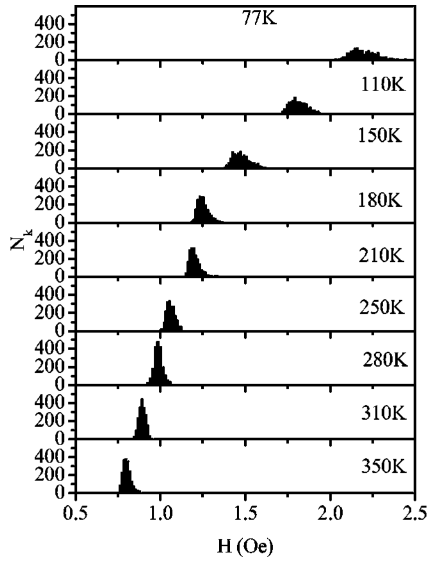

FIG. 2. Switching field distribution at different temperatures.

two pinning mechanisms of the closure domain wall. First, there is a magnetoelastic pinning of the closure domain structure coming from the mechanical stresses induced during the preparation process of the microwires together with the additional stresses induced by the glassy coating, especially as a consequence of their different thermal expansion coefficients. The corresponding temperature dependence of the coercivity can be expressed in the form

$$H_\sigma(T) \propto cM_s^x(T)(1 + r(\Delta T))^{1/2}. \quad (16)$$

Here c=const σ_r, x depends on the origin of the magnetostriction, $x \in \langle 1;1.5\rangle$ (for Co-based alloys, x=1), and $r \approx E(\alpha_g - \alpha_m)/\sigma_r$. σ_r are the stresses induced during the microwire production, E is Young's modulus and α_g and α_m are the thermal expansion coefficients of the glass coating and the metallic nucleus, respectively.

Second, the pinning of the closure domain on the atomic level defects was found to be responsible for the increase of the switching field at low temperatures. The coercivity which rises from this mechanism can be expressed as

$$Hp(T) \propto \frac{1}{M_s}\frac{\varepsilon_p^2\rho_p}{kT}F(T,t), \quad (17)$$

where ε_p corresponds to the interaction energy of the mobile defects with the local spontaneous magnetization, ρ_p is the density of the mobile defects, k is a Boltzmann constant and $F(T,t)$ is a relaxation function:[18] $F(T,t)=(1-e^{(-t/\tau)})$, 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 preexponential factor and Q denotes to the activation energy of the mobile defects.

As observed in Fig. 3, the temperature dependence of the switching field distribution can be satisfactorily fitted to the sum of the contributions of these two mechanisms which have different temperature dependencies:

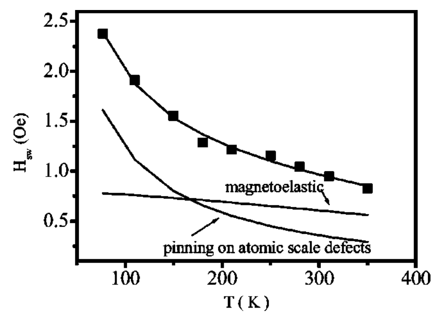

FIG. 3. The temperature dependence of the main switching field. The full lines correspond to the fits to Eqs. (16)–(18).

$$H_{sw}(T) = cM_s^x(T)(1 + r(\Delta t))^{1/2} + nF(T,t)/(M_sT), \quad (18)$$

where x=1 for a Co based sample, and $n \approx (\varepsilon_p^2\rho_p/k)$ is proportional to the number of defects when $F(T,t)$ is assumed to be a monotonously increasing function of the measuring temperature at the measured times and temperatures.

B. Switching field distribution shape

As it rises out from the measuring method, the probability of the closure domain to overcome the energy barrier in some field interval $(H;H_{max})$ depends on the fact that the barrier was not overcome at lower fields $(0;H)$. Such a problem arises purely from the adopted method and is not related in any way to the underlying physics of the effect.

To determine properly the switching field distribution the distributions given in Fig. 2 can be recalculated using a similar procedure given in Ref. 16. An approximated expression can be calculated relating the probability density given in Eq. (12) as

$$\frac{dp}{d(\Delta h)} = \left[\sum_{k=1}^{k_m}\frac{N_k}{\sum_{i=k}^{k_m} N_i}\right]^{-1}\frac{N_k}{\delta h\sum_{i=k}^{k_m} N_i}, \quad (19)$$

where $\delta_h=(H_{max}-H_{min})/k_m$ and k_m is a number of divisions of the interval $(H_{min}-H_{max})$. Equation (19) takes into account the fact that conditions for appearance of large and small fluctuations are nonequivalent due to the experimental method. Moreover it gives us the possibility to confront the theory with the experiment by comparing Eq. (12) with Eq. (19).

Figure 4 presents the dependence of the probability density to overcome the energy barrier $(dp/d(\Delta h))$ on the reduced magnetic field (Δh). In opposition to Eq. (12) which gives a linear dependence of $\ln(dp/d(\Delta h))$ on $(\Delta h)^{3/2}$, a more complex behavior is found. Anyway, it can be understood in terms of the thermoactivated model taking into account the above considered two domain wall potentials.

As it was shown elsewhere,[9] the coercivity mechanism which comes from the pinning of the domain wall has two origins in bistable amorphous wires: the magnetoelastic pinning at stress centers and the pinning at the defects on an atomic scale. These two interactions have different potentials. The magnetoelastic pinning comes from a long-range

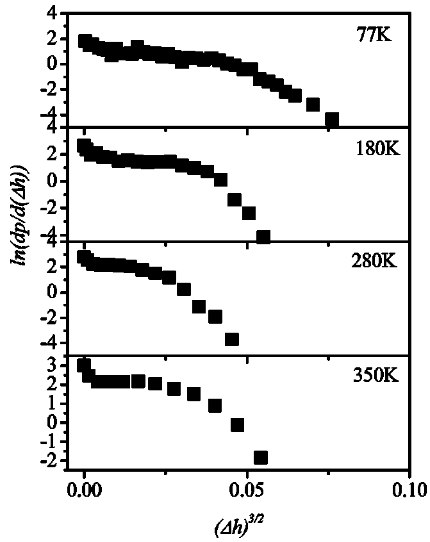

FIG. 4. Dependence of the probability density on the reduced magnetic field.

interaction between stress applied on the microwires and stress centers that are present due to its amorphous nature. Its energy density is proportional to the magnetostriction, λ_s, through the equation

$$E_\sigma = (3/2)\lambda_s \sigma, \tag{20}$$

and due to its long-range nature, the shape of the magnetoelastic potential is wide (see Fig. 5).

On the other hand, the potential coming from the short-range pinning of the domain wall on atomic scale defects is narrow and its width is comparable to the domain wall. It rises from the structural relaxation of the mobile atomic scale defects which lose their mobility when the temperature decreases. Frozen defects stabilize the local magnetic moments within the domain wall as well as within the wall. The stabilization energy of the domain wall with the defects can be analytically given by[10]

$$E_p = (2/15)\varepsilon_p^2(\rho_0/kT)[-2\delta_0 + x^2/\delta_0]F(T,t), \tag{21}$$

where δ_0 is the domain wall thickness and x is displacement from the domain wall equilibrium position.

Surely there is another contribution to the long range potential coming from the magnetostatic energy which rises due to the decrease of the stray field energy by the formation of the closure domain structure as well as the shape anisotropy energy. But, as it was deduced from the temperature dependence of the switching field, their contributions can be neglected in comparison to the two above mentioned contributions.

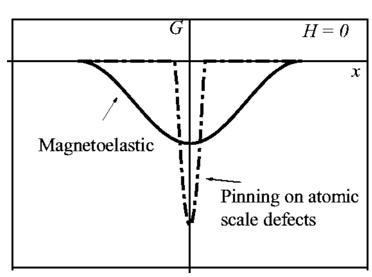

FIG. 5. Closure domain wall potentials coming from the magnetoelastic pinning and pinning on the atomic scale defects.

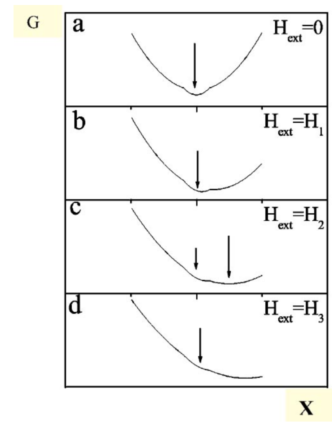

FIG. 6. Schematic dependence of the total free energy G of the closure domain wall under the action of external magnetic field $H_{ext}(H_1 < H_2 < H_3)$ on the domain wall position x.

Contrary to the temperature dependence of the switching field, the decomposition of the linear dependence of the $\ln(dp/d(\Delta h))$ on $(\Delta h)^{3/2}$ into two linear dependencies corresponding to two potentials is not so easy a job. The below described construction approaches the change of the domain wall free energy potentials under the action of the external field.

So, the total free energy of the closure domain wall without the action of an applied magnetic field is given as a sum of these two potential as schematically shown in Fig. 6(a). In the low magnetic field regime, H_1, the free energy minimum is given by the shape of the potential coming from the short range pinning [Fig. 6(b)]. At an intermediate field, H_1, two local minima co-exist in the free energy of the domain wall [Fig. 6(c)]. This happens for specific shapes of the two potentials. For a field close to the switching field, the second minimum disappears, but small local minimum still exist coming from the pinning on the atomic scale defects [Fig. 6(d)].

Based on this assumption, the specific shape of the probability density dependence on the reduced magnetic field can be explained. It has no linear dependence, but it consists of the addition of two linear dependencies as schematically depicted in Fig. 7. The slope of the linear curve is given at each temperature [according to Eq. (12)] by the susceptibility and coercivity (through the α) of the corresponding potentials.

Considering the thermally activated model,[15,17] the switching field distribution width should be a function of $T^{2/3}$. Nevertheless, the width of the distribution experimentally decreases with the temperature which can be explained considering a temperature dependence of the switching field. Its value can be determined from the relationship[13]

$$\Delta H = (1/\alpha)^{2/3} C H_{sm}, \tag{22}$$

where Q is a constant parameter. In the first approximation, the distribution width is proportional to the coercivity (see

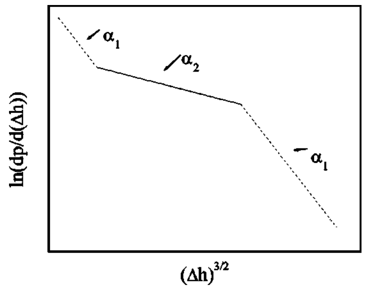

FIG. 7. Theoretical dependence of the probability density on the reduced magnetic field.

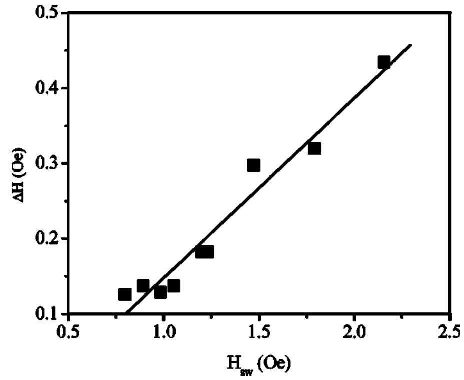

FIG. 8. The dependence of the switching field distribution width ΔH_{sw} on the switching field H_{sw} at different temperatures.

Fig. 8), where the temperature dependence of α has been neglected.

V. CONCLUSION

The specific experimental temperature dependence of the switching field distribution in bistable amorphous wires has been interpreted in terms of the thermoactivated model. The model is extended in the sense that two terms are found to contribute to the total potential wall responsible for the switching mechanism of these microwires: the long range magnetoelastic and short range pinning of the domain wall on the atomic level defects. The total free energy is then given by the addition of the two contributions which results in the specific shape of the domain wall potential. Under the action of the external magnetic field, the minimum of the free energy is given by a different parts of the potential which results in the specific switching field distribution.

The distribution width is found to be proportional to the switching field at different temperature in opposition to the previous results.

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