A model is studied for the theoretical description of nanoscale magnetic films with high perpendicular anisotropy. In the model the magnetic film is described in terms of single domain magnetic grains with Ising-like behavior, interacting via exchange as well as via dipolar forces. Additionally, the model contains an energy barrier and a coupling to an external magnetic field. Disorder is taken into account in order to describe realistic domain and domain wall structures. The influence of a finite temperature as well as the dynamics can be modeled by a Monte Carlo simulation.

Many of the experimental findings can be investigated and at least partly understood by the model introduced above. For thin films the magnetisation reversal is driven by domain wall motion. The results for the field and temperature dependence of the domain wall velocity suggest that for thin films hysteresis can be described as a depinning transition of the domain walls rounded by thermal activation for finite temperatures.

I. INTRODUCTION

We focus on thin ferromagnetic films with high perpendicular anisotropy like CoPt and MnBi which can be used as magneto-optical storage media. In these films two different mechanisms can be thought to dominate the reversal process: either nucleation or domain wall motion \[^{[4]}\]. Which of these mechanisms dominates a reversal process depends on the interplay of the different interaction forces between domains with different magnetic orientation. In a recent experiment on Co\(_{28}\)Pt\(_{72}\) alloy films \[^{[5]}\] a crossover from magnetisation reversal dominated by domain growth to a reversal dominated by domain wall motion \[^{[5]}\] has been achieved by simulations of a micromagnetic model \[^{[5]}\] the film is thought to consist of cells on a square lattice with a square base of size \(L \times L\) where \(L = 200\)\(\text{Å}\) and height \(h\) of 100\(\text{Å}\). Due to the high anisotropy of the Co\(_{28}\)Pt\(_{72}\) alloy film the grains are thought to be magnetised perpendicular to the film only with a uniform magnetisation \(M_s\) which is set to the experimental value of \(M_s = 365\)\(\text{kA/m}\) for the saturation magnetisation in these systems \[^{[6]}\]. The grains interact via domain wall energy and dipole interaction. The coupling of the magnetisation to an external magnetic field \(H\) is taken into account as well as an energy barrier which has to be overcome during the reversal process of a single cell.

From these considerations it follows that the change of energy caused by reversal of a cell \(i\) with magnetisation \(L^2hM_s\sigma_i\) with \(\sigma_i = \pm 1\) is:

\[
\Delta E_i = -\frac{1}{2}LhS_w\Delta\sigma_i \sum_{<j>} \sigma_j + \frac{\mu_0}{4\pi}M_s^2Lh^2\Delta\sigma_i \sum_j v(\sigma_j, r_{ij}) - \mu_0 H L^2 h M_s \Delta \sigma_i
\]  

(1)

The first term describes the wall energy \(\Delta E_w\). The sum is over the four next neighbors and for \(S_w\) we use a value of \(S_w = 0.0022\)\(\text{J/m}^2\) which is approximately 50% of the Bloch-wall energy \(S_B\) for this system \[^{[2]}\]. The reason for this reduction of the grain interaction energy compared to the Bloch-wall energy is, that the crystalline structure of the system is interrupted at the grain boundary and also that due to their irregular shape the grains are not connected via their complete surface \(Lh\).

In the second term describing the dipole coupling \(\Delta E_d\) the sum is over all cells. \(r_{ij}\) is the distance between two cells \(i\) and \(j\) in units of the lattice constant \(L\). For large distances it is \(v(\sigma_j, r_{ij}) = \frac{2}{r_{ij}}\). For shorter distances a more complicated form which is a better approximation for the shape of the cells which we consider can be determined numerically and was taken into account.

The third term describes the coupling \(\Delta E_H\) to an external field \(H\).

Additionally, an energy barrier \(\delta_i\) must be considered which describes the fact that a certain energy is needed to reverse an isolated cell by domain wall motion through the grain (see also \[^{[2]}\]). We assume that during the reversal process the energy barrier has its maximum value \(LhS_w\) when the domain wall is in the center of the cell, i.e. when half of the cell is already reversed. Consequently,
the energy barrier which is relevant for the reversal process is reduced to 
\[ \delta = \max(0, L_hS_h - \frac{1}{2}((\Delta E_w + \Delta E_d + \Delta E_H))) \] . The simulations are in good qualitative agreement with experiments using 
\[ S_h = 0.0007 J/m^2 \].

In order to simulate the CoPt films realistically disordered has to be considered. Obviously, the grain sizes are randomly distributed. In the model above this corresponds to a random distribution of \( L \) which can hardly be simulated exactly since it modulates the normalized cell distance \( r_{ij} \) of the dipole interaction. Therefore, as a simplified ansatz to simulate the influence of disorder we randomly distribute \( L \) in the energy term that describes the coupling to the external field. Here a random fluctuation of \( L \) is most relevant, since this term is the only one scaling quadratic ly with \( L \). In the simulations we use a distribution which is Gaussian with width \( \Delta \). Note, that through this kind of disorder our model is mapped on a random-field model.

The simulation of the model above was done as in an earlier publication via Monte Carlo methods using the Metropolis algorithm with an additional energy barrier. Since the algorithm satisfies detailed balance and Glauber dynamics it allows the investigation of thermal properties as well as the investigation of the dynamics of the system. For temperatures \( T \to 0 \) the Monte Carlo algorithm passes into a simple energy minimization algorithm with single spin flip dynamics, so that also the case of zero temperature can be investigated.

The size of the lattice was typically \( 150 \times 150 \). The dipole interaction was taken into account without any cut-off or mean field approximation.

III. HYSTERESIS AND DYNAMICS

Fig. 1 shows a simulated hysteresis loop. The loop is nearly rectangular. Here, the reversal is dominated by domain wall motion. Once a nucleus begins to grow the domain wall motion does not stop until the magnetisation has completely changed. The nucleation field for the simulation is too high compared to the corresponding experimental results. There are several possible reasons for this effect. First, the nucleation field which in the case of domain wall motion dominated reversal is the field where domain wall motion starts depends on the size and on the shape of the nucleus. In an experimental situation a nucleus can be very large e. g. a scratch while there is no such artificial nucleus in our simulation except of the disorder which leads only to nuclei of very small size. However, it is not the aim of our simulations to calculate the nucleation field accurately. Rather, the simulations are thought to contribute to a better understanding of the fundamental properties of the system.

Since the reversal is driven by domain wall motion, the velocity of the domain wall is a central quantity which we will investigate in the following. For the determination of the domain wall velocity within the simulation we start with a system that has a nucleus of circular shape with a radius of 19 cells in the center of the \( 150 \times 150 \) system. When we switch on the driving field, from the nucleus a domain starts to grow. For the better observation of the domain growth, in our flip-algorithm we do not consider cells that are not connected to the growing domain, i.e. we exclude the possibility of additional nucleation. Otherwise we have – at least for finite temperatures – the problem that spontaneously new nuclei are build by thermal activation which with increasing radius overlap with the original domain. Fig. 3 shows - as an example - the \( v(H) \) behavior from the simulations for \( T = 0K \) and different fields. For \( r > 75 \) the domain reaches the boundary of the system and consequently - \( v(H) \) saturates. For zero temperature the domain wall is pinned for lower fields, i.e. after a short period of rearrangement of the domain wall the domain wall movement stops and the radius remains constant. The pinning of the domain wall is due to energy barriers which follow from the disorder, the dipole field, and the intrinsic energy barrier of the single cell. For finite temperatures, the domain wall velocity is always finite.

For \( 20 < r < 75 \) the slope of the \( v(H) \) curve is approximately constant and \( v \) can be determined by fitting to a straight line. Fig. 3 shows the dependence of the domain wall velocity on the driving field for \( T = 0K \) and different fields. For \( r > 75 \) the domain reaches the boundary of the system and consequently - \( v(H) \) saturates. For zero temperature there is a sharp depinning transition at a critical field \( H_c \) from a pinned phase with \( v = 0 \) to a phase with finite domain wall velocity. This transition can be interpreted in terms of a phase transition with \( \nu \to (H - H_c)^\theta \) for \( H > H_c \) where in our case the critical exponent is \( \theta \approx 1 \), a value which is the mean field result for a moving elastic interface in a random field. Also, this value has been observed in simulations of a soft spin model with random-fields. For finite temperatures the transition is smeared since for finite temperatures there is even for \( H < H_c \) for each energy barrier a finite probability that the barrier can be overcome by thermal fluctuations. The corresponding waiting time can be expected to be exponentially large so that for \( H < H_c \) the domain wall velocity should decrease like \( \ln v \sim (H - H_c) / T \). To illustrate this in Fig. 5 we show the corresponding semi-logarithmic scaling plot. As we expect, the data for the two different finite temperatures collapse for \( H < H_c \) on a straight line. For \( H > H_c \) thermal activation is obviously less relevant. Here, the dynamics is dominated by the zero temperature depinning transition.

To conclude, the results for the domain wall velocity suggest that for zero temperature the hysteresis driven by domain wall motion can be understood as a depinning transition of the domain walls. For finite temperature the transition is rounded and for fields smaller than the depinning field the domain wall movement is dominated by thermal activation. A paper on a comparison of these theoretical results with experimental measure-
ments of the domain wall velocity in CoPt alloy films is in preparation.

ACKNOWLEDGMENTS

The author thanks K. D. Usadel for helpful discussions and for critically reading the manuscript.

[1] J. Pommier, P. Meyer, G. Penissard, J. Ferre, P. Bruno, and D. Renard, Phys. Rev. Lett. 65, 2054 (1990)
[2] T. Kleinefeld, J. Valentin, D. Weller, JMMM 148, 249 (1994)
[3] J. Valentin, T. Kleinefeld, D. Weller, J. Phys. D 29, 1111 (1996)
[4] U. Nowak, IEEE Trans. Mag. 31, 4169 (1995)
[5] W. Andra, H. Danan, and R. Mattheis, Phys. Stat. Sol. (a), 125, 9 (1991)
[6] J. Harzer, RWTH Aachen, Ergebnisbericht (1992)
[7] U. Nowak, U. Ruediger, P. Fumagalli, G. Guntherodt, Phys. Rev. B 54, 13017 (1996)
[8] K. Binder, Monte Carlo Methods in Statistical Physics (Springer-Verlag, Berlin 1979)
[9] M. Kardar and D. Ertas, in Scale Invariance, Interfaces, and Non-Equilibrium Dynamics, edited by A. McKane, M. Droz, J. Vannimenus, and D. Wolf, Plenum Press, New York 1995, pa. 89
[10] H. Leschhorn, J. Magn. Magn. Mat. 104-107, 309 (1992)
[11] K. D. Usadel and M. Jost, J. Phys. A26, 1783 (1993)

FIG. 1. Hysteresis loop for a 100Å film; $T = 300K$

FIG. 2. Domain configuration of a $150 \times 150$ system during the reversal; $T = 300K$

FIG. 3. Radius of the domain versus time for the same fields as in Fig. 3 $T = 0$; solid lines are best fitted.

FIG. 4. Domain wall velocity versus driving field for $T = 0, 300$, and 600K; solid lines are guides to the eye.

FIG. 5. Scaling plot from Fig. 4