Thermal stability of coupled ferromagnetic and superparamagnetic particles

Vladimir L. Safonov and H. Neal Bertram
Center for Magnetic Recording Research, University of California - San Diego, 9500 Gilman Dr., La Jolla, CA 92093-0401, U.S.A.

March 22, 2022

We consider a single-domain ferromagnetic particle with uniaxial anisotropy coupled to a single-domain soft ferromagnetic particle (superparamagnetic particle). The problem of thermally agitated magnetization reversal in this case can be reduced to the random magnetization dynamics of the first particle with an effectively larger anisotropy field. The magnetic external field is also altered in a manner that depends on the sign of the coupling and can be either enhanced or suppressed.

I. INTRODUCTION

Thermal stability is of great importance in ultra-high density magnetic recording. The principal element of magnetic memory is a single-domain ferromagnetic particle. The magnetic moment of this particle is agitated by thermal fluctuations and can reverse to the opposite direction. In the absence of inter-particle interactions the magnetization reversal time is determined by the Néel-Arrhenius exponent with a definite energy barrier (e.g., [1], [2], [3]). The problem for the mean magnetization reversal time has been rigorously solved in the framework of a new theoretical approach, in which the coherent rotation of the magnetization is analyzed in terms of the motion of a damped nonlinear oscillator subject to random Gaussian forces (thermal agitation) [4], [5], [6].

Recently it has been experimentally demonstrated that the thermal stability of magnetic memory can increase if the ferromagnetic particles from the recording layer are antiferromagnetically coupled to an auxiliary layer of smaller ferromagnetic particles (AFC media) [7], [8]. Simple theoretical explanations of this phenomenon are given, for example, in Refs. [4], [9]. The basic model consists of two single-domain hard particles (recording and stabilizing) coupled by an antiferromagnetic exchange interaction.

Here we analyze the possibilities of this two-particle model to enhance thermal stability of recording media. We consider the thermally agitated magnetization reversal, when a hard ferromagnetic particle interacts with a soft ferromagnetic (superparamagnetic) particle. We use the new theoretical approach for coupled damped nonlinear oscillators subject to random thermal agitation.

II. MODEL

Let us consider a ferromagnetic particle with magnetization \( M_1(\mathbf{M}_1 = M_{s,1}) \) and volume \( V_1 \) and a superparamagnetic particle with \( M_2 \) \((\mathbf{M}_2 = M_{s,2}) \) and \( V_2 \), respectively (Fig.1). To describe the coherent motion of the particle magnetizations we introduce the classical spins \( S_j = M_j V_j / \hbar \gamma \) \((j = 1, 2)\), where \( \gamma \) is the gyromagnetic ratio and \( \hbar \) is Plank’s constant.

The energy of the combined system includes uniaxial anisotropy of the ferromagnetic particle along \( z \) axis, Zeeman energy, exchange and magnetostatic interparticle interactions:

\[
\mathcal{E} = -K_u V_1 \left( \frac{S_1^z}{S_1} \right)^2 - \hbar \gamma H_0 \cdot (S_1 + S_2) + JS_1 \cdot S_2 + (\hbar \gamma)^2 \left[ \frac{S_1 \cdot S_2}{r^3} - \frac{3 (S_1 \cdot r) (S_2 \cdot r)}{r^5} \right].
\]

Here \( K_u \) is the uniaxial anisotropy constant, \( H_0 \) is the external magnetic field and \( J \) is the exchange integral \((J > 0 \text{ corresponds to antiferromagnetic and } J < 0 \text{ to the ferromagnetic coupling, respectively})\). The last term in \( \mathcal{E} \) describes the magnetostatic interactions (dipole-dipole approximation), where \( r \) is the vector between particle centers.

For simplicity we shall confine our analysis to two particle configurations applicable for the longitudinal (Fig.1a) and perpendicular (Fig.1b) recording media with magnetic fields parallel to each uniaxial anisotropy axis. Thus Eq. (1) becomes

\[
\mathcal{E} = -K_u V_1 \left( \frac{S_1^z}{S_1} \right)^2 - \hbar \gamma H_0 (S_1^z + S_2^z) + \hbar \gamma (H_{int} S_1^z S_2^z + H_{\perp,z} S_1^z S_2^y + H_{\perp,y} S_1^y S_2^z).
\]

Here \( H_{\perp,y} = J / \hbar \gamma + \hbar \gamma / r^3 \). For configurations shown in Fig.1 the interaction fields \( H_{int} \) and \( H_{\perp,x} \) are defined by

\[
H_{int}^{(a)} = \frac{J}{\hbar \gamma} + \frac{\hbar \gamma}{r^3},
\]

\[
H_{\perp,x}^{(a)} = \frac{J}{\hbar \gamma} - 2 \frac{\hbar \gamma}{r^3}.
\]

*E-mail: vsafonov@ucsd.edu
for Fig. 1a and
\[ H_{\text{int}}^{(b)} = \frac{J}{\hbar \gamma} - \frac{\hbar \gamma}{r^3}, \]
\[ H_{\perp,z}^{(b)} = \frac{J}{\hbar \gamma} + \frac{\hbar \gamma}{r^3} \]
for Fig. 1b.

### III. Dynamic Equations

In order to describe spin (magnetization) dynamics, we introduce complex variables \( a_j^+ \) and \( a_j \) (classical analog of creation and annihilation operators) of nonlinear oscillators defined by the Holstein-Primakoff transformation \[3]:

\[
S_j^x = S_j - N_j, \quad N_j = a_j^+ a_j,
\]
\[
S_j^y = \frac{a_j + a_j^+}{2} \sqrt{2S_j - N_j},
\]
\[
S_j^z = \frac{a_j - a_j^+}{2i} \sqrt{2S_j - N_j}.
\]

This transformation gives a convenient representation of spins as oscillators in the vicinity of equilibrium \( S_1 = (0, 0, S_1) \) and \( S_2 = (0, 0, S_2) \).

The dynamic equations for \( a_1 \) and \( a_2 \), defined by \( \frac{da_j}{dt} = -i \partial (E/h)/\partial a_j^+ \), can be written as

\[
\frac{d}{dt} a_1 = -i \gamma H_{\text{eff},1} a_1 - i \gamma F_{12},
\]
\[
\frac{d}{dt} a_2 = -i \gamma H_{\text{eff},2} a_2 - i \gamma F_{21}.
\]

Here

\[
H_{\text{eff},1} = H_0 + H_K S_1^z - H_{\text{int}} S_2^z - \frac{J}{\hbar \gamma} - \frac{\hbar \gamma}{r^3},
\]
\[
H_{\text{eff},2} = H_0 - H_{\text{int}} S_1^z - \frac{J}{\hbar \gamma} + \frac{\hbar \gamma}{r^3}
\]
are the effective fields, \( H_K = 2K_u/M_s \) is the anisotropy field and

\[
F_{mn} = H_{\perp, n} \frac{\partial S_m^x}{\partial a_n^+} S_m^x + H_{\perp, n} \frac{\partial S_m^y}{\partial a_n^+} S_m^y.
\]

The principal oscillatory motion of \( a_1(t) \) and \( a_2(t) \) are determined by their effective fields: \( a_1 \propto \exp(-i \gamma H_{\text{eff},1} t) \) and \( a_2 \propto \exp(-i \gamma H_{\text{eff},2} t) \). Because, in general, \( H_{\text{eff},1} \neq H_{\text{eff},2} \), the fast oscillating terms \( F_{12} \) and \( F_{21} \) in \[1\] vanish. The magnetization fluctuations of the superparamagnetic particle occurs at a much faster rate \( \sim f_2 \) than the reversal rate of the ferromagnetic particle. \( f_2 \) is an attempt frequency of the second particle (a negligibly small energy barrier is assumed). Thus, the longitudinal component \( S_1^z \) is the slowest variable in the system, and the analysis of the coupled system can be reduced to the stochastic dynamics of the first particle in an averaged field from the second particle. This field appears as a feedback of the superparamagnetic particle on the instantaneous ordered state \( S_1^z \) of the first particle. The corresponding effective energy is equal to

\[
\mathcal{E}_2 \simeq -\hbar \gamma H_{\text{eff},2} S_2^z.
\]

In this mean field approximation \[10\], it follows that for times \( \gg f_2^{-1} \) an averaged thermal spin polarization of the second (superparamagnetic) particle \( \langle S_2^z \rangle \) is:

\[
\langle S_2^z \rangle = S_2 B_{S} \left( \frac{S_2 \hbar \gamma H_{\text{eff},2}}{k_B T} \right),
\]

where

\[
B_S(Sx) = 1 - \frac{1}{S} \left( \frac{1}{\exp(x) - 1} - \frac{2S + 1}{\exp[(2S + 1)x] - 1} \right)
\]

is the Brillouin function. This averaged spin polarization \( \langle S_2^z \rangle \) induces an effective field (or, response field) back to the first particle. As a result the effective energy of the first particle can be written as

\[
\mathcal{E}_1 \simeq -\hbar \gamma \frac{H_K}{2S_1} \langle S_1^z \rangle^2 - \hbar \gamma (H_0 - H_{\text{int}} \langle S_2^z \rangle) S_1^z.
\]

From \[3\] and \[1\] one can see that \( H_{\text{int}} \) is enhanced for antiferromagnetic \( J > 0 \) exchange in the longitudinal case (Fig. 1a) and for ferromagnetic exchange in the perpendicular case (Fig. 1b).

### IV. Approximate Solution

Let us consider a small polarization of the superparamagnetic particle. In this case Eq. \[11\] is simplified to

\[
\langle S_2^z \rangle \simeq \frac{S_2 (S_2 + 1) \hbar \gamma H_{\text{eff},2}}{3 k_B T}
\]

and the effective energy \[13\] becomes

\[
\mathcal{E}_1 \simeq -\hbar \gamma \frac{\tilde{H}_K}{2S_1} \langle S_1^z \rangle^2 - \hbar \gamma H_{\text{eff},1} S_1^z.
\]

Here

\[
\tilde{H}_K = H_K \left( 1 + 2S_1 \frac{H_{\text{int}}}{H_K} \right)
\]

is the effective anisotropy field and

\[
H_{\text{eff},1} = H_0 (1 - \zeta)
\]

is the effective external field on the first particle,

\[
\zeta = \frac{S_2 (S_2 + 1) \hbar \gamma H_{\text{int}}}{3 k_B T}.
\]
From (16) and (18) we see that the effective anisotropy field $H_{\text{eff}}$ always increases ($H_{\text{int}} \zeta \propto H_{\text{int}}^2$) for any $H_{\text{int}} \neq 0$. On the contrary, the effective external magnetic field $H_{\text{eff},1}$ (17) depends on the sign of $H_{\text{int}}$: it can be both enhanced ($\zeta \propto H_{\text{int}} < 0$) and suppressed ($\zeta > 0$). The feedback effect from the superparamagnetic particle quadratically depends on the saturation magnetization and volume of this particle: $\zeta \propto S^2_{V} \propto M^2_{V} V^2_{V}$. This quadratic dependence also means that the corresponding feedback effect of two superparamagnetic particles of volume $V_{2}/2$ is at least two times smaller. Thus it is more effective to use a single larger superparamagnetic particle than a collection of smaller ones.

The energy (14) is an exact analog to the case of random motion of one single-domain particle when the external magnetic field is parallel to the anisotropy axis. Thus, we can use exact analytic formulas (15), (16), (17) for the mean first passage time obtained for this case. It should be noted that this approach makes it possible to describe experimental data for the reversal field versus pulse time over 13 orders from nanoseconds to hours (6), including both purely dynamic and thermally assisted reversal.

V. RESULTS AND DISCUSSION

Let us estimate Eqs. (13) and (17). For the ferromagnetic particle we take the following parameters at room temperature: $K_{v}V_{l}/k_{B}T = 60$, $H_{K} = 5$ kOe and $M_{s,1} = 150$ emu/cc. This means $V_{1} \simeq 6.62 \cdot 10^{-18}$ cm$^3$ and $S_{1} \simeq 5.36 \cdot 10^4$. Taking $S_{2} \simeq 700$ and $H_{\text{int}} = 0.3$ Oe, one obtains: $\zeta \simeq 0.022$ and $2S_{1}\zeta H_{\text{int}}/H_{K} \simeq 0.14$. Thus, we see that the effective anisotropy field is relatively small in the effective magnetic field $H_{\text{eff}}\simeq 0.978H_{0}$ and larger in the effective anisotropy field $H_{K} = 1.14H_{K}$. For the opposite sign of interaction (ferromagnetic exchange between particles) $H_{\text{int}} = -0.3$ Oe and with the same parameters, one has $H_{\text{eff},1} \simeq 1.022H_{0}$ and $H_{K} = 1.14H_{K}$.

In Fig. 2 reversal field versus pulse time is shown for the above two cases. We also show the corresponding curve for one ferromagnetic particle ($H_{\text{int}} = 0$) and Sharrock’s (Néel-Arrhenius) dependence. The thick solid line and broad dashed curve represent antiferromagnetic and ferromagnetic coupling, respectively. The slight difference is due to the larger effective field and hence reduced stability of the ferromagnetic case. In the thermal region ($H < H_{K}$) there is more curvature than in the Néel-Arrhenius approximation. This is because the nonlinear oscillator model contains a field dependent attempt frequency. Note that formulas (17) are valid for long time scale $f^{-1} \sim 10^{-10} - 10^{-9}$ sec. In the dynamic reversal region particles seems to reverse independently.

ACKNOWLEDGMENT

This work was partly supported by matching funds from the Center for Magnetic Recording Research at the University of California - San Diego and CMRR incorporated sponsor accounts.

[1] W. F. Brown, Jr., Phys. Rev. 130, 1677 (1963).
[2] D. Weller and A. Moser, IEEE Trans. Magn. 35, 4423 (1999).
[3] H. J. Richter, J. Phys. D: Appl. Phys. 32, R147 (1999).
[4] V. L. Safonov, J. Magn. Magn. Mater. 195, 526 (1999); J. Appl. Phys. 85, 4370 (1999).
[5] V. L. Safonov and H. N. Bertram, J. Appl. Phys. 87, 5681 (2000).
[6] H. N. Bertram and V. L. Safonov, IEEE Trans. Magn. 36, 2447 (2000).
[7] V. L. Safonov and H. N. Bertram, Magnetization dynamics and thermal fluctuations in fine grains and films. pp.81-109 in: ”The Physics of Ultra-High-Density Magnetic Recording”, Eds. M. Plumer, J. van Ek and D. Weller (Springer-Verlag, 2001).
[8] X. Wang, H. N. Bertram, and V. L. Safonov, J. Appl. Phys. 92, 2064 (2002).
[9] E. N. Abarra, A. Inomata, H. Sato, I. Okamoto, and Y. Mizoshita, Appl. Phys. Lett. 77, 2581, (2000).
[10] E. E. Fullerton, D. T. Margulies, M. E. Schabes, M. Carey, B. Gurney, A. Moser, M. Best, G. Zeltzer, K. Rubin, H. Rosen, and M. Doerner, Appl. Phys. Lett. 77, 3806 (2000).
[11] M. E. Schabes, E. E. Fullerton, and D. T. Margulies, IEEE Trans. Magn. 37, 1432 (2001).
[12] H. J. Richter and Er. Girt, IEEE Trans. Magn. 37, 1441 (2001); H. J. Richter, Er. Girt, and H. Zhou, Appl. Phys. Lett. 80, 2529 (2002).
[13] T. Holstein and H. Primakoff, Phys. Rev. 58, 1098 (1940).
[14] A. Aharoni, Introduction to the theory of ferromagnetism (Clarendon press, Oxford, 1996).
[15] M. P. Sharrock, J. Appl. Phys. 76, 6413 (1994).

Figure captions

Fig.1. Particle configurations for a) longitudinal and b) perpendicular recording. The larger volume represents the ferromagnetic particle.

Fig.2. Reversal field versus pulse time.
This figure "FSFig1.jpg" is available in "jpg" format from:

http://arxiv.org/ps/cond-mat/0212041v1
This figure "FSFig2.jpg" is available in "jpg" format from:

http://arxiv.org/ps/cond-mat/0212041v1