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# Heat-driven spin torques in antiferromagnets

Marcin Białek<sup>✉</sup>, Sylvain Bréchet and Jean-Philippe Ansermet<sup>✉</sup>

Institute of Physics, École Polytechnique Fédérale de Lausanne (EPFL), Lausanne, Switzerland

E-mail: [marcin.bialek@epfl.ch](mailto:marcin.bialek@epfl.ch)

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## Abstract

Heat-driven magnetization damping, which is a linear function of a temperature gradient, is predicted in antiferromagnets by considering the sublattice dynamics subjected to a heat-driven spin torque. This points to the possibility of achieving spin torque oscillator behavior. The model is based on the magnetic Seebeck effect acting on sublattices which are exchange coupled. The heat-driven spin torque is estimated and the feasibility of detecting this effect is discussed.

Keywords: antiferromagnet, spin caloritronics, spin torque oscillator, spin current, antiferromagnetic resonance


## 1. Introduction

Antiferromagnets present magnetization dynamics that fall naturally in the THz range. This is of interest for spintronics, since such high frequencies are expected to become of use, given the ever increasing amount of data that is needed to be transferred [1, 2]. While some components and devices are now available to fill the so-called THz gap [3, 4], much remains to be developed, in particular when it comes to integrated circuits operating at such high frequencies. Therefore, there has been much interest in exploring ways in which spintronics strategies can make use of the naturally high-frequency dynamics of antiferromagnets. Since the overwhelming majority of antiferromagnets are insulators, we examine in this paper the possibility that spin torque can be generated by heat rather than a charge current. Thermal spin torque adds a new damping term to the equations defining the high frequency dynamics of spins, which is the prerequisite for making a heat-driven spin torque oscillator.

Spin torque oscillators are magnetic nanostructures in which a large current drives the precession of the magnetization. The earliest studies of spin transfer torque consisted of detecting sharp features in the differential conductance of multilayers [5] or spin valves [6], due to the precession

of the magnetization [7, 8]. The magnetic field at which these oscillations occurred could be provided by the addition of a magnetic pillar to a spin valve [9]. Oscillations were obtained in spin valves with a perpendicular-anisotropy material at zero applied field and low threshold currents [10]. It has been demonstrated that oscillations can be induced in an insulating magnet by driving a DC current through a Pt layer deposited on it [11]. The spin orbit torque is responsible for driving the magnetization into sustained precession, an effect that was first shown in a metallic ferromagnet [12]. Since these spin torque oscillators are nanostructures, their output power is very low. However, it has been shown that coupled oscillators can be phase-locked [13, 14] with a spin Hall magnetoresistance effect as one of the coupling mechanisms [15, 16]. High-coherence emission was achieved by entraining the oscillator with a driving microwave signal [17, 18].

The spin torque generated by heat currents was demonstrated to drive magnetization dynamics in an insulating ferromagnet [19, 20] and to induce a magnetization flop [21]. The heat current flowing through the interface of a ferromagnet and a nonmagnetic metal was demonstrated to produce pure spin currents [22] and affect the magnon properties in ferromagnets [23–25]. It was predicted theoretically that a thermal current in conducting ferromagnets should produce measurable spin torques [26] and vice versa [27]. Thermal spin torques explain the activation of domain wall motion in insulating ferromagnets [28]. It has been suggested that thermal spin torques might act on skyrmions in an insulating chiral magnet [29].

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The interest in high frequency devices motivates research on oscillators based on spin transfer torques in antiferromagnets [30]. It has been pointed out that the spin Hall effect may excite the magnetization of a canted antiferromagnet [31]. A spin current can be injected into an antiferromagnetic insulator without having a charge current, and it may be possible to drive an oscillation of its magnetization [32]. The scheme in that case consists of driving a large charge current in a Pt layer adjacent to an antiferromagnet. The spin current thusly produced is injected into an antiferromagnet, which must have biaxial anisotropy in order for the spin current to couple to the magnetization sublattices. For practical applications, once a sustained precession is achieved, it must be converted into an electrical signal. The proof of principle for this has been established by showing that antiferromagnetic resonances can be electrically detected using the inverse spin Hall effect [33].

A promising material for a high frequency oscillator might be a half-metallic ferrimagnet [34, 35] in which the two antiparallel magnetization sublattices are compensated, so that demagnetization fields are negligible and the spin transfer torque is optimal [36]. It is very difficult to grow such materials, but success was reported in growing  $\text{Mn}_2\text{Ru}_{0.33}\text{Ga}$  with the Curie temperature of 550 K [37]. Similarly to antiferromagnets, ferrimagnets present very fast magnetization dynamics without external magnetic fields [38]. Ferrimagnetic  $\text{Mn}_3\text{Ge}$  was demonstrated to produce THz emission when exposed to ultra-short laser pulses [39, 40], as were earlier canted antiferromagnets [41, 42].

In this paper, we focus on the potential use of a heat-driven spin torque in antiferromagnets. First, we review how magnetization dynamics in a ferromagnet is affected by a temperature gradient. In light of this result, we analyze how heat might affect magnetization dynamics in an antiferromagnet by considering it as composed of two opposing magnetization sub-lattices that are exchange-coupled. We conclude by discussing which materials might be appropriate for detecting a heat-driven spin torque. The mechanism discussed in this paper is clearly distinct from spin reorientation effects that are induced by intense optical pumping [43, 44]. Such ultrafast reorientation can excite the magnetic modes of antiferromagnets [45]. Contrary to these experiments, here we consider small temperature changes and the effect of a steady temperature gradient. Under such conditions, the spin Seebeck effect in Pt-antiferromagnet interfaces subject to magnetic fields was reported in  $\text{Cr}_2\text{O}_3$  [46] and  $\text{MnF}_2$  [47]. In this paper, we are considering bulk effects without external magnetic fields.

## 2. Ferromagnet in a temperature gradient

When a ferromagnet is subjected to a temperature gradient, the magnetization field is non-uniform because it depends on temperature. It is possible to obtain an equation for magnetization dynamics with the method of Landau and Lifshitz, which is based on a variation principle [48], assuming that the energy is a function not only of  $\mathbf{M}$ , but of  $\nabla \times \mathbf{M}$  also [49]. This gives an extra term to the Landau–Lifshitz–Gilbert equation, which now reads:

$$\dot{\mathbf{M}} = \gamma \mathbf{M} \times \mathbf{B}_{\text{eff}} + \frac{\alpha}{M_s} \mathbf{M} \times \dot{\mathbf{M}} + \gamma \mathbf{M} \times \mathbf{B}_{\text{TST}}, \quad (1)$$

where  $\mathbf{B}_{\text{eff}}$  is the usual effective field (applied field  $\mathbf{B}_0$ , demagnetization and anisotropy fields),  $\gamma$  is the gyromagnetic ratio,  $\alpha$  is a damping parameter and  $M_s$  is the modulus of the magnetization. The last term is a heat-induced spin torque. It is cast in such a way that it can be thought of as resulting from a heat-induced field  $\mathbf{B}_{\text{TST}}$  [20], with

$$\mathbf{B}_{\text{TST}} = -\mu_0 \left( \mathbf{k}_T \cdot \nabla^{-1} \right) \mathbf{M}. \quad (2)$$

Here, a thermal wave vector  $\mathbf{k}_T$  is introduced in order to clarify the units of the new term that are added to the Landau–Lifshitz–Gilbert equation, with

$$\mathbf{k}_T = \frac{-1}{\chi^2} \frac{d\chi}{dT} \nabla T. \quad (3)$$

The linear response deduced from the modified Landau–Lifshitz–Gilbert equation (equation (1)) reads

$$\delta \dot{\mathbf{M}} = \gamma (\delta \mathbf{M} \times \mathbf{B}_{\text{eff}}) + \frac{\alpha}{M_s} \mathbf{M}_s \times \delta \dot{\mathbf{M}} - \gamma \mu_0 \mathbf{M}_s \times \left( \mathbf{k}_T \cdot \nabla^{-1} \right) \delta \mathbf{M}. \quad (4)$$

## 3. Antiferromagnet in a temperature gradient

We use the Pincus approach to describe magnetization dynamics in an antiferromagnet [50], where each sublattice magnetization experiences the combination of the applied field  $\mathbf{B}_0$ , the anisotropy field  $\mathbf{B}_A$  and an exchange field proportional to the magnetization of the other sublattice. The proportionality constant is noted  $\lambda$ . Including a damping term, this leads to the following equation of evolution for the magnetization of each sublattice:

$$\frac{d\mathbf{M}_1}{dt} = \gamma \mathbf{M}_1 \times (\mathbf{B} + \lambda \mathbf{M}_2) + \frac{\alpha}{M_s} \mathbf{M}_1 \times \frac{d\mathbf{M}_1}{dt}, \quad (5)$$

$$\frac{d\mathbf{M}_2}{dt} = \gamma \mathbf{M}_2 \times (\mathbf{B} + \lambda \mathbf{M}_1) + \frac{\alpha}{M_s} \mathbf{M}_2 \times \frac{d\mathbf{M}_2}{dt}. \quad (6)$$

Therefore, we can consider in a first approximation that under a heat current, each sublattice follows a dynamical equation similar to equation (4), with the applied field comprising an anisotropy field  $\mathbf{B}_A$  and an exchange field  $\mathbf{B}_i$  given by  $\lambda \mathbf{M}_i$ , ( $i = 2, 1$ ). Thus, we can use equation (4) to obtain linearized equations for the antiferromagnet in a temperature gradient

$$\begin{aligned} \delta \dot{\mathbf{M}}_1 &= \gamma (\delta \mathbf{M}_1 \times (\mathbf{B}_A + \lambda \mathbf{M}_2)) + \gamma (\mathbf{M}_1 \times \lambda \delta \mathbf{M}_2) \\ &+ \frac{\alpha}{M_s} \mathbf{M}_1 \times \delta \dot{\mathbf{M}}_1 - \gamma \mu_0 \mathbf{M}_1 \times \left( \mathbf{k}_T \cdot \nabla^{-1} \right) \delta \mathbf{M}_1, \end{aligned} \quad (7)$$

$$\begin{aligned} \delta \dot{\mathbf{M}}_2 &= \gamma (\delta \mathbf{M}_2 \times (\mathbf{B}_A + \lambda \mathbf{M}_1)) + \gamma (\mathbf{M}_2 \times \lambda \delta \mathbf{M}_1) \\ &+ \frac{\alpha}{M_s} \mathbf{M}_2 \times \delta \dot{\mathbf{M}}_2 - \gamma \mu_0 \mathbf{M}_2 \times \left( \mathbf{k}_T \cdot \nabla^{-1} \right) \delta \mathbf{M}_2. \end{aligned} \quad (8)$$

The equilibrium state, in the absence of an applied field, is given by

**Table 1.** Material values for the estimation of a temperature gradient  $\nabla T$  required to obtain a 100% change in the AFMR linewidth ( $\Delta\alpha = \alpha$ ).

Material	$T$ (K)	$T_N$ (K)	$\omega^{-1}d\omega/dT$ (K <sup>-1</sup> )	$M_s^{-1}dM_s/dT$ (K <sup>-1</sup> )	$\nabla T$ (K cm <sup>-1</sup> )
MnF <sub>2</sub>	$T < T_N$	60	$\omega^{-1}d\omega/dT \approx M_s^{-1}dM_s/dT$ [54]	—	—
DyFeO <sub>3</sub>	$\approx 50$	640	$1 \cdot 10^{-2}$ [55]	$-6 \cdot 10^{-4}$	$2 \cdot 10^2$
DyFeO <sub>3</sub>	$\approx 300$	640	0 [55]	$-2 \cdot 10^{-3}$	$1 \cdot 10^3$
Cr <sub>2</sub> O <sub>3</sub>	$\approx 200$	300	0 [56]	$-3 \cdot 10^{-3}$	$7 \cdot 10^2$

$$\mathbf{M}_1 = -\mathbf{M}_2 = \mathbf{M}_s. \quad (9)$$

The thermal effect is a small perturbation of the normal modes,  $m_{\pm} = \delta M_{1x} \pm \delta M_{1y} + \delta M_{2x} \pm \delta M_{2y}$ , with the linear response to an excitation  $h_{\pm}$  given by [48]

$$m_{\pm} = \chi_{af} h_{\pm} = \frac{2\gamma^2 \mu_0 M_s B_A}{(\omega_+ - \omega)(\omega_- - \omega) + 2i\alpha\omega\gamma B_E} h_{\pm}, \quad (10)$$

where  $B_E = \lambda M_s$ . Thus, the resonance line width of these modes is given by [48]

$$\Delta B_{\pm} = \frac{2\alpha\omega_{\pm} B_E}{\gamma B_A}. \quad (11)$$

We can quantify how the heat-driven spin torque affects the magnetization dynamics as a change to the damping parameter  $\alpha$ . Hence, the relaxation terms in equations (7) and (8) can be rewritten as

$$\delta \dot{\mathbf{M}}_i|_{\text{rel}} = \frac{\alpha}{M_s} \mathbf{M}_i \times \left( \frac{d}{dt} - \frac{\gamma\mu_0 M_s}{\alpha} \mathbf{k}_T \cdot \nabla^{-1} \right) \delta \mathbf{M}_i \quad (i = 1, 2). \quad (12)$$

If  $\mathbf{M}_i$  contains a term of the form  $\mathbf{M}_0 \exp(i\mathbf{k} \cdot \mathbf{r} - i\omega t)$  and if the temperature gradient is in the direction of  $\mathbf{k}$ , then the effective damping parameter is given by [20]

$$\alpha_{\text{eff}} = \alpha + \Delta\alpha = \alpha + \frac{\gamma\mu_0 M_s k_T}{\omega k}. \quad (13)$$

As demonstrated experimentally in [20], the longer the wavelength of the magnetization wave, the larger the heat-driven spin torque. This is expressed in equation (13) by the  $1/k$  factor. It is analogous to the dependence on  $1/q$  of the thermoelectric effect, where  $q$  is the electric charge of the charge carrier. In a metal, the thermoelectric coefficient  $\varepsilon$  can be approximated by [51]

$$\varepsilon \approx \frac{k_B T}{e T_F}, \quad (14)$$

where  $T_F$  is the Fermi energy,  $T$  temperature and  $e$  the electron charge. On a microscopic level, the spatial variation of thermal energy  $k_B T$  is counterbalanced by the spatial variation of electrostatic energy,  $eV$  where  $V$  is the electrostatic voltage. Likewise, in the magnetic Seebeck effect [52], the spatial variation of thermal energy is counterbalanced by the spatial variation of the magnetic energy, expressed as  $\mathbf{M} \nabla B$ .

#### 4. Heat-driven self-oscillations in an antiferromagnet

Equation (13) is essentially the same as for ferromagnets. However, the expression for  $k_T$  is different for antiferromagnets.

It can be calculated using equations (3) and (10), the approximate Kittel formula for the antiferromagnetic resonance  $\omega \approx \gamma \sqrt{2B_E B_A}$  and under the assumption that  $\omega_+ = \omega_- = \omega$ :

$$k_T = \frac{2\alpha\gamma\lambda^2 M_s}{\mu_0 \omega} \left( \frac{1}{M_s} \frac{dM_s}{dT} - \frac{1}{\omega} \frac{d\omega}{dT} \right) \nabla T. \quad (15)$$

We can estimate the order of magnitude of the contribution from the thermal spin torque to the width of the resonance. The relative change of line width due to the thermal spin torque is given by

$$\frac{\Delta\alpha}{\alpha} = \frac{\gamma\mu_0 M_s k_T}{\alpha\omega k} = \frac{B_E}{B_A} \left( \frac{1}{\omega} \frac{d\omega}{dT} - \frac{1}{M_s} \frac{dM_s}{dT} \right) \frac{\nabla T}{k}. \quad (16)$$

The form of equation (16) shows that the effect of the temperature gradient is present if the dependence of the antiferromagnetic resonance frequency  $\omega$  is different from that of the sublattice magnetization  $M_s$ . This difference reflects the role of the anisotropy in these materials, since the resonance frequency is proportional to  $\sqrt{B_A}$ .

The sublattice magnetization  $M_s$  is strongly temperature-dependent in antiferromagnets with a virtually flat dependence around  $T = 0$  K and a very fast change near the Néel temperature, where it can be estimated using the power law

$$M_s(T) = D M_s(0) \left( 1 - \frac{T}{T_N} \right)^{\beta}, \quad (17)$$

with  $D \approx 1.1$ . Mössbauer measurements of the temperature dependence of the internal field showed that  $\beta = 1/3$  in orthoferrites [53]. Hence,  $dM_s/dT$  is of the order of  $M_s T_N^{-1}$  at a temperature  $T \approx \frac{2}{3} T_N$ . We can expect a similar power law for other materials. If we assumed the value of  $1/2$  for  $\beta$  as predicted by the Landau theory of phase transitions, or the  $2/3$  Bloch law, we would get the same order of magnitude for the estimate of this coefficient. In the material specific estimations that follow, we assume that  $B_E/B_A \approx 10^2$  and the wave vector  $k \approx 2 \cdot 10^4 \text{ m}^{-1}$  which is that of an electromagnetic radiation at 300 GHz in a material with a dielectric constant of  $\sim 9$ . The obtained values of  $\nabla T$  required for a self oscillation are presented in table 1.

In some antiferromagnetic compounds, like MnF<sub>2</sub>,  $\omega$  is almost proportional to  $M_s$  [54], thus making  $k_T = 0$ . However, in some antiferromagnetic materials, which show nontrivial dependences of  $\omega$  on temperature, the thermal spin torque effect on the AFMR width might not be negligible. For example, the canted antiferromagnet DyFeO<sub>3</sub> with  $T_N = 640$  K experiences a spin reorientation phase transition at  $T_M \approx 45$  K owing to a strong change in the magnetic anisotropy [55]. Because of this phenomenon, one of the antiferromagnetic

resonances softens around  $T_M$ , and above this temperature its frequency rises rapidly with  $\omega^{-1}d\omega/dT \approx 10^{-2} \text{ K}^{-1}$ . In this case, a gradient of just  $20 \text{ K cm}^{-1}$  would be sufficient to achieve a 10% change in the resonance width (equation (16)).

However, a linewidth change around  $T_M$  might be very difficult to detect because of the strong temperature dependence of the resonance. Thus, for the purpose of an experimental demonstration of the effect, materials with  $\omega(T)$  showing an extremum, so that  $d\omega/dT = 0$ , might be more promising. The frequency of the AFMR in  $\text{DyFeO}_3$  reaches a maximum at around room temperature. There, a gradient of  $100 \text{ K cm}^{-1}$  would be sufficient to achieve a 10% change in the resonance width. Smaller values of  $\nabla T$  might be needed in materials with lower  $T_N$  and where an extremum of the AFMR frequency is closer to  $T_N$ . For example, the AFMR in a colinear antiferromagnet  $\text{Cr}_2\text{O}_3$  has a maximum at  $T = \frac{2}{3}T_N \approx 200 \text{ K}$  [56]. In this case,  $\nabla T \approx 70 \text{ K cm}^{-1}$  would be sufficient to observe a 10% change in the resonance width (equation (16)). This particular material has the additional advantage of presenting a magnetoelectric coupling [57].

Values of  $\nabla T$  required to obtain self oscillation behavior (table 1) are out of reach using bulk materials at temperatures where  $d\omega/dT = 0$ . However, it has been shown that such large temperature gradients can be achieved in thin films using a laser beam [58]. Joule heating in nanostructures was also shown to give rise to such large temperature gradients [22]. Estimated gradient values are an order of magnitude smaller in the case of antiferromagnets with the spin reorientation transition, like  $\text{DyFeO}_3$  at around 50 K.

## 5. Summary



We have used a variational approach to identify a heat-driven spin torque when an antiferromagnet is subjected to a temperature gradient. The key parameter that characterized this effect is the temperature dependence of the sublattice magnetization and the anisotropy field. Temperature dependence of the magnetization is typically described in terms of thermal excitation of magnons. Hence, this variational approach has some connection with other theoretical approaches that consider magnon statistics [59–61]. We have applied the concept of a heat-driven spin torque to examine how it modifies the magnetization dynamics in an antiferromagnet. We used a simple description which consists of assuming that each sub-lattice of the antiferromagnet is exchanged-coupled to the other one. Finally, we carried out a quantitative estimate of the effect of a heat-driven spin torque on the width of the resonance, and found that its strength depends critically on the difference between the temperature dependences of the resonance and the magnetization. We suggest antiferromagnets with a spin reorientation transition as candidates where this effect might be observed.

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## ORCID iDs

Marcin Białek  <https://orcid.org/0000-0002-3209-3489>  
Jean-Philippe Ansermet  <https://orcid.org/0000-0002-1307-4864>

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