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Chapter 6

Nonlocal spin Seebeck effect in the bulk easy-plane antiferromagnet NiO

6.1 Abstract

We report the observation of magnon spin currents generated by the spin Seebeck effect (SSE) in a bulk single crystal of the easy-plane antiferromagnet NiO. A magnetic field induces a non degeneracy and thereby an imbalance in the population of magnon modes with opposite spins. A temperature gradient then gives rise to a nonzero magnon spin current. This SSE is measured in both a local and a nonlocal geometry at 5 K in bulk NiO. The magnetic field dependence of the obtained signal is modeled by magnetic field splitting of the low-energy magnon modes, affecting the spin Seebeck coefficient. The relevant magnon modes at this temperature are linked to cubic anisotropy and magnetic dipole-dipole interactions. The nonlocal signal deviates from the expected quadratic Joule heating by saturating at a current from around 75 \( \mu \text{A} \) in the injector. The magnon chemical potential does not decay exponentially with distance and inhomogeneities may be the result of local magnon accumulations.

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6.2 Introduction

Magnon spintronics is a field where spin currents are carried by magnons that exist in tunable magnets for information processing [1]. Generation of spin currents in magnets is feasible by using the spin Hall effect (SHE) in a metal injector strip creating the magnons which travel through the magnetic material to be subsequently detected at a second detecting strip [2]. Antiferromagnets (AFMs) do not possess stray fields and can therefore be exploited over a wide range of parameters such as external magnetic field and device size. Recently, this nonlocal technique has effectively been employed for the uniaxial AFMs $\alpha$-Fe$_2$O$_3$ [3], MnPS$_3$ [4], and Cr$_2$O$_3$ [5]. Despite their potential for spin transport by both magnons and spin superfluidity [6], this geometry has not been employed for easy-plane antiferromagnets like NiO.

Magnons are quasi-particles carrying spin angular momentum which enables the transfer of spins in (insulating) magnets as waves of spin rotations of the magnetic moments. An easy-axis antiferromagnet has left-handed ($\alpha$) and right-handed ($\beta$) magnon modes in which energies are equal but spins are opposite. Magnon interconversion is expected to equal the respective magnon chemical potentials $\mu^\alpha_m = \mu^\beta_m = \mu_m$, the deviation from the equilibrium magnon population. Magnon injection then creates a finite $\mu_m$ which drives the transport of magnon spins, following the regular discussion for magnon transport [7]. This description is equivalent to that in Refs. [8, 9], where $\mu^\alpha$ and $\mu^\beta$ are regarded as equal but opposite as resulting in opposite spin currents. Magnon spins can be injected at the interface with a paramagnetic heavy metal using the SHE or in the bulk magnet.

In the first method, $\alpha$ ($\beta$) magnons are created ( annihilated) if the accumulated spin direction at the interface is parallel to an $\alpha$-magnon spin resulting in an increase (decrease) in the magnon chemical potential $\mu_m$. In the latter method, heating by the injector sets up a thermal magnon current of both modes, which diffuses from the hot region to the cold region, the spin Seebeck effect (SSE). In a ferromagnet (FM) this leads to finite spin currents even without an applied magnetic field since there is one kind of magnon whose polarity is determined solely by the magnetization. In AFMs, however, there is no inherent population imbalance between the modes when they are degenerate. These modes carry equal but opposite spin currents and no net spin current arises due to a thermal gradient, so there is no SSE in the absence of a magnetic field. The degeneracy of these magnon modes is lifted by a magnetic field, creating an imbalance in their population and thereby net magnon spin currents can be created [10, 11].

The different modes are coupled via inelastic magnon-magnon scattering, al-
lowing for energy interchange between different modes and magnon relaxation. This results in a slight suppression of the spin Seebeck coefficient but will largely leave the transport and its magnon conductivities of both magnon modes intact \[12\]. The magnon depletion is expected to decay exponentially with distance in bulk magnets\[12\] \[13\]. Accumulation of magnons at interfaces can be observed as a sign change in $\mu_m$ \[13\].

In this chapter, the nonlocal observation of the spin Seebeck effect in an easy-plane AFM shows that there is an imbalance between populations of the magnon modes with opposite spins. This results in a net spin current driven by the temperature gradient. The transport of magnons does not require such an imbalance. However, no electrically injected magnon spin currents, which rely on a significant spin transmission by an exchange interaction between the Pt and the NiO, have been observed. The SSE-generated spin currents have been generated at 5 K in bulk NiO containing multiple domains. There is no need of arranging the magnetic domains by the exchange interaction with a FM layer, in contrast to Ref. \[14\] where a FM seed layer is required to obtain a SSE signal from a 200-nm-thick NiO layer. The SSE amplitude as a function of the magnetic field strength is modeled by magnon mode splitting, creating the imbalance in the magnon population of the modes. The magnon chemical potential shows some local variation and shows an increase in noise by increasing the distance from the injector.

### 6.3 Theory

#### 6.3.1 Magnon modes

Easy-axis antiferromagnets have magnon modes that are typically in the terahertz regime. Easy-plane AFMs, however, have a more complex magnon dispersion, which is extended to lower energies. When considering only the exchange and anisotropy, a gap appears between the two modes \[14\]. Although NiO has a simple rock salt structure, magnetic dipole-dipole interactions and cubic anisotropy can give rise to multiple low-energetic precession modes, the dispersions of which depend on the magnetic field \[16\]. At low temperatures (a few Kelvins), the gap from Zeeman splitting is of the order of the thermal energy and therefore could induce a nonzero spin Seebeck coefficient.

A magnetic field influences the spin current via the dispersion consisting of multiple modes with different magnon spin polarizations. Reference \[15\] treats the dis-
Figure 6.1: (a) Device structure and probes on the NiO bulk crystal with the spin injector (left Pt strip) and detector (right Pt strip) at distances \(d\) apart, ranging from 250\,nm to 7\,\mu m. An in-plane magnetic field \(B\) is applied with a clockwise angle \(\alpha\) with the \(y\)-axis. A 100-\(\mu\)A current is sent through the injector leading to Joule heating and a radial heat gradient \(J_q\) indicated by the red arrows. A surplus of magnons on the hot side flow to the cold side, leaving behind a negative \(\mu_m\). At interfaces, magnons accumulate and can contribute to \(\mu_m\) as observed in thin films of yttrium iron garnet (YIG) on gadolinium gallium garnet (GGG), the distribution of which is reproduced here \cite{13}. \(\mu_m\) is normalized in the scale. A finite \(\mu_m\) results in spin transport between the Pt and the NiO via the spin mixing conductance. The inverse spin Hall effect consequently causes a voltage locally, \(V_l\), and nonlocally, \(V_{nl}\). (b) The magnon dispersion [Eq. (6.1)] after Ref. \cite{15} of two modes as a function of a magnetic field along the easy plane considering exchange and Zeeman interaction. The inset shows the full range of the wave number \(q\). (c) Magnon energy at the \(q = 0\) point after Ref. \cite{16} as a function of a magnetic field along [110] when taking magnetic dipole-dipole interactions and cubic anisotropy into account. This breaks the symmetry and splits the magnon energies with [2\(\bar{1}\)1] spin directions. The energies are given in Kelvin (black) and meV (red).

persion of the \(\alpha\) and \(\beta\) modes having an offset at the zero \(q\) point even without an applied magnetic field. The offset is said to arise from the hard-axis anisotropy and is further influenced by the Zeeman interaction when applying a magnetic field. The
dispersion is then given by

\[
\omega_{\alpha,\beta}^2 \gamma_k^2 = (H_e + H_{ep} + \frac{H_{ha}}{2})^2 - \frac{1}{4} H_{ha}^2 + H^2 - H_e^2 \gamma_k^2 \\
\pm \left( H_e^2 H_{ha} \gamma_k^2 + 4H^2((H_e + H_{ep} + \frac{H_{ha}}{2})^2 - H_e^2 \gamma_k^2) \right)^{\frac{1}{2}}
\]

(6.1)

where \( \gamma_k = \cos(\frac{1}{2}k a_l) \), \( H_e=968.4 \text{ T} \), \( H_{ha}=635 \text{ mT} \), and \( H_{ep} = 11 \text{ mT} \) are the structure factor, the exchange interaction and the hard-axis and easy-plane anisotropy, respectively \([15, 17–19]\). The offset causes an unequal population of these modes, especially at temperatures lower than the offset temperature of the \( \alpha \) mode. Since their magnon spin directions are considered opposite \([10]\), this results in a net magnon spin required for a nonzero SSE. Figure 6.1(b) shows that the offset in the dispersions as a function of magnetic field further increases, enlarging the net magnon spin.

However, when additionally considering the symmetry breaking magnetic dipole-dipole interactions and the cubic anisotropy as done by Milano and Grimsditch \([16]\), multiple low-energy modes appear which are shown in Fig. 6.1(c) for different domains and the respective magnetic moment directions as a function of the magnetic field strength. Energetically higher modes are not considered since the measurements are performed at 5 K. Under the influence of a magnetic field within a \(< 111> \) easy plane (along \([110]\)), the \([2\overline{1}1]\) magnon modes and to a lesser extend the \([11\overline{2}]\) magnon modes split, the \([1\overline{1}2]\) magnon modes remain degenerate, while the \([2\overline{1}1]\) modes are soft and become unstable from 0.55 T. A magnetic field thus causes an imbalance in the occupation of these modes which have different magnon spin direction. This leads to an imbalance in the magnon populations and opens up the opportunity to investigate these magnons with the SSE.

### 6.3.2 Spin Seebeck effect

When there is a net magnon spin population in a magnetic material, a temperature gradient can drive a magnon spin current \( J_s = -(\sigma_m \nabla \mu_m + S \nabla T) \) via the SSE. The spin Seebeck coefficient \( S \) has a field and a temperature dependence. The flow of magnons creates a negative magnon chemical potential \( \mu_m \) near the injector. Boundary conditions at interfaces \([13]\), and, possibly, domain walls and defects, lead to magnon accumulation and reflection resulting in a positive sign of \( \mu_m \) at a distance from the injector. Shown in Fig. 6.1(a) is the distribution of \( \mu_m \) as a result of such reflections at the interface of a thin film of YIG on GGG. A spin current enters a Pt strip via a finite spin mixing conductance where it is converted to a charge current.
by the inverse spin Hall effect.

A SSE generated spin current has been observed from NiO when grown on a FM to force the preference of one type of the many possible domains by the exchange interaction which is stronger with an uncompensated ferromagnetic <111> layer. These domains are said to have an anisotropy-induced splitting without a magnetic field and having only one type of domain would result in a nonzero SSE \[14\]. Without the seed layer of Py below the grown NiO, no SSE was observed as the spin currents originating from different domains would cancel \[14\]. Signals in such systems show hysteresis and decrease in size when decreasing the temperature, vanishing below 100 K \[20\].

### 6.3.3 Spin Hall magnetoresistance

The spin Hall magnetoresistance (SMR) is the first harmonic response and can be obtained simultaneously with the second harmonic SSE with a lock-in technique \[21\]. The SMR of the Pt injector strip is sensitive to the magnetic moments underneath it, even to the Néel vector in antiferromagnets \[22\]. The SHE deflects electrons in a direction depending on their electron spin, resulting in the accumulation of electron spins at the interface with NiO. The direction of these electron spins is affected by the interaction with the magnetic moments in the NiO via the spin transfer torque. This exchange interaction is maximal when the directions of the magnetic moments and the accumulated electron spin are perpendicular. The electron spin is reflected back into the Pt and subsequently deflected by the inverse spin Hall effect determined by the electron spin. Absorption of the spin by the magnet thereby affects the path traveled by the electrons and influences the longitudinal resistivity \(\rho_L\) of the Pt layer by \[22\]

\[
\rho_L = \rho + \Delta \rho_0 + \Delta \rho_1 < 1 - n_x^2 >
\]

where \(< n_x >\) is the average of the Néel vector along \(\hat{x}\) just below the Pt injector. This technique thus indicates the influence of the magnetic field on the magnetic moments; i.e., it gives information about the magnetic order and domain wall growth.

### 6.4 Properties of NiO

NiO is a cubic material with antiferromagnetic interaction due to superexchange between two Ni atoms via an oxygen ion. Together with magnetic dipole-dipole
interactions and a small cubic anisotropy this results in the spins aligning in ferromagnetic \{111\} planes which are intercoupled antiferromagnetically [23]. The magnetic moments themselves align along [112] within these \{111\} planes, although a minor diversion can be induced by anisotropy [23] or rhombohedral distortion. Due to magnetostriction the crystal is rhombohedrally distorted along the <111> directions and magnetic twin (T) domains are formed [24]. Within a T domain, the three easy axes give rise to corresponding spin rotation (S) domains. By introducing a magnetic field, the degeneracy of energetically equivalent domains is lifted, resulting in a redistribution of these domains by movement of the domain walls. The direct influence on the spin rotation causes movement of the S domains. Domain walls can influence the rotation of the Néel vector and thereby both the SMR and the SSE.

6.5 Methods

The bulk NiO sample was commercially obtained and polished along a <111> plane as described in Ref. [22]. Thereafter, the devices were fabricated using electron beam lithography. No etching was performed before the sputtering of the 5-nm-thick, 20-µm-long, and 100-nm-wide Pt strips. Three devices were fabricated with distances between the Pt strips of 250 nm to 7 µm. The electrically and thermally originated signals are measured by obtaining the 1\textsuperscript{st} and 2\textsuperscript{nd} harmonic separately but simultaneously with the lock-in technique [21].

6.6 Results

The SMR measurements show a $\sin^2\alpha$ angle dependence [see Fig. 6.2(a)]. The changes in the magnetic moments are such that they tend to align perpendicular to the magnetic field direction, so as to maximize the negative Zeeman energy. The signals are therefore 90° angular shifted as compared to Pt|FM systems, confirming antiferromagnetic order in the material [22]. Similar to our earlier work on this sample [22], the amplitudes of such rotation measurements initially increase quadratically as a function of the field strength and a saturation sets in which is established around 6 T [see Fig. 6.2(b)]. This field dependence is believed to be originating from both anisotropy and domain wall movement. As domains with the magnetic easy plane in the <111> surface plane become more favoured over other domains, these domains grow in size [22, 25]. The domain size at small field must be smaller than the Pt strips size to follow the same field dependence as a Hall bar device. This
Figure 6.2: (a) Spin Hall magnetoresistance as a function of the in-plane rotation angle $\alpha$, performed at 6.3 T and 5 K. $R_\alpha$ is the angle-dependent resistance of the Pt strip and $R_0$ is the base resistance of 5.17 k$\Omega$. (b) Amplitude of the SMR signals as a function of the magnetic field strength, showing a similar curve as reported in [22] for both devices. (c) The locally observed signal changes from the spin Seebeck effect as a function of the magnetic field angle at 6.3 T. (d) The amplitude of the signal increases monotonely with field with little offset, modeled by the three methods described in the main text. The error bars in panels (b) and (d) represent the fit uncertainty.

agrees well with the observed domain size of $<1 \mu$m [26]. A domain wall can affect magnon spin currents and therefore the distribution of magnon chemical potential as well. At saturated field strengths, the magnetic moments are coherently rotated by the magnetic field and the crystal is more or less in a single magnetic domain.

The locally measured current-induced SSE shows an angular dependence and signal size which are similar to those of local measurements of Pt on FMs [see Fig. 6.2(c)]. The noise of 40 $V/A^2$ is relatively large in comparison to Pt|FM systems and might be originating from domain walls that move due to the changing magnetic field direction. There is a background signal of which the origin can be other heat-related effects such as the spin Nernst or the Righi-Leduc effect. The size of the
Figure 6.3: Local SSE using a transverse Hall bar geometry. The signals are obtained at a magnetic field strength of 8 T at 5 K and 33 K. Due to the larger surface area, the effects of domains are diminished and this graph validates the usage of a sinusoidal fit for the data.

The SSE amplitude as a function of the magnetic field strength shows on average a monotonous increase for all three devices as shown in Fig. 6.2(d) apart from the substantial variation between measurements.

Local SSE signals have been measured by a Hall bar geometry as well [see Fig. 6.3]. The range of the rotation angle $\alpha$ is larger in this graph, validating the sinusoidal fit used in Fig. 6.2(c). Further, it is shown that at higher temperatures the signal vanishes. Up to room temperature, the measurements did not show a local or nonlocal SSE signal within the noise.

The SSE amplitude as a function of the magnetic field strength was fitted using the contribution of the different modes by methods A and B described in Ref. [15] and method C from Ref. [27]. The magnon dispersion from Eq. (6.1) has been used in method B and the zero $q$-point value has been altered by the data of Ref. [16] presented in Fig. 6.1(c) for methods A and C, assuming that only small $q$-values play a role at the low temperature of 5 K. It is assumed that the magnons are close to the thermodynamic equilibrium, allowing the use the Bose-Einstein distribution for both magnon branches at a temperature of the phonon bath [12, 15]. For methods A and B, the spin Seebeck coefficient of the splitted magnon modes is given by

$$S_S^z = S_0 \int dkk^2 \left\{ \frac{\hbar \omega_{\beta k}}{e^{\beta k B} - 1} \omega_{\beta k} \nu_{\beta k}^2 - \frac{\hbar \omega_{\alpha k}}{e^{\beta k B} - 1} \omega_{\alpha k} \nu_{\alpha k}^2 \right\}$$  \hspace{1cm} (6.3)

where $S_0 = \frac{\hbar}{6\pi^2 k_B T^2}$, $\omega_{\mu k}$ is the field dependent frequency and $\eta_{\mu k}$ is the magnon relaxation rate of the magnon mode $\mu$ as function of the wave vector $k$. Further we need the magnon group velocity $\nu_{\mu k} = \partial \omega_{\mu k} / \partial k$ and the relaxation rate $\eta_\mu = 1 + (1240q + 5860q^3)(T/300)^5$ [27]. The values are numerically solved at a large set
Figure 6.4: The nonlocally obtained SSE signals. (a) The angle-dependent resistance at 5 K and 6.3 T. The angular dependence is similar to that of Pt on FMs and is opposite to that of the local signal. The background resistance is small in comparison to the local background resistance. (b) The second harmonic has a negative sign at distances smaller than 500 nm, similar to in thin ferromagnetic films, but after the sign change the signal increases monotonously with distance. The size of the error bars increase with distance and have a large variation between devices. (c) The signal size increases in a comparable fashion with increasing magnetic field strength as the local signal. (d) The SSE does not increase quadratically as a function of the current. Instead, the signal shows some current threshold behavior until 25 $\mu$A and increases with field until it is saturated around 75 $\mu$A.

of field strengths whose interpolating function can be fitted to the SSE amplitude data as a function of the field strength. Using the method described in Ref. [15] this amounts to $S_S^z = 7.8 \times 10^{-11}$ erg cm$^{-1}$ K$^{-1}$ for the dispersion in Eq. (6.1) (method A) and $S_S^z = 2.4 \times 10^{-11}$ erg cm$^{-1}$ K$^{-1}$ using the adjusted dispersion (method B) at 7 T.

Both the near linear increase and the relatively small offset in the data are not represented by method B which instead shows a large offset and an approximately quadratic increase with field strength. Possibly these modes disappear due to spin reorientations when applying a field. By following the model described in Ref. [15]...
and using Eq. (6.1), the field dependence shows saturating behavior and a less significant offset. When using the dispersion including the effects of magnetic dipole-dipole interactions and cubic anisotropy, the SSE field dependence resembles the data and no offset is present. At 20 K the signal has reduced to $(3.2 \pm 0.3) \times 10^{-3}$ V A$^{-2}$ at 1.5 T, in agreement with this explanation.

The temperature gradient is calculated using its relation with the signal size given by [28]

$$V_{SSE} = R_N l \lambda_N \frac{2e}{h} \theta_{SH} \tanh \left( \frac{t_N}{2\lambda_N} \right) S_S^2 \nabla T,$$

(6.4)

where $R_N = 5.17 \, \text{k}\Omega$, $l = 20 \, \mu\text{m}$, $t_N = 8 \, \text{nm}$, $\lambda_N = 1.1 \, \text{nm}$, and $\theta_{SH} = 0.08$ are the Pt bar resistance, the length, the thickness, the spin diffusion length and the spin Hall angle, respectively [29]. Further, it is assumed that the NiO thickness $\gg$ the relaxation length. The temperature gradient along $\hat{z}$ near the injector with a current of $100 \, \mu\text{A}$ is calculated to be $2.50 \times 10^6$ K m$^{-1}$ for method C, method A gives $7.51 \times 10^5$ K cm$^{-1}$, and method B results in $2.30 \times 10^3$ K cm$^{-1}$. This is lower compared to the calculated average temperature gradient resulting from a similar geometry on YIG at 300 K of $1.6 \times 10^8$ K cm$^{-1}$ [13], which indicates an overestimation of the calculated $S_S^2$ value. On the other hand, the thermal conductivity can be different in NiO at 5 K as compared to that in YIG at 300 K.

Figure 6.4(a) shows the angular dependence of a SSE signal obtained nonlocally at the detector and is similar to that of devices on thick YIG with a strip distance in the same range [13]. The distance dependence of the signal, shown in Fig. 6.4(b), shows a sign change around $d \approx 500$ nm, indicating that $\mu_m$ turns positive. Moreover, after the sign change, the signal seems to increase with increasing distance. In a YIG thin film on a GGG substrate, such sign changes are subscribed to boundary conditions at the interface of the thin film and the paramagnetic substrate. The interface will conduct little magnons while the heat is transported into GGG. Therefore, the magnons accumulate and are reflected causing $\mu_m$ to turn positive at a certain distance from the injector and $\mu_m$ goes to zero at large distances [13]. With further increasing distances the signal in FM's drops according to a diffusion-relaxation model [2]. However, single-domain bulk FMs lack these boundaries and no positive $\mu_m$ is observed [8]. Recently, it is shown that the rotation of the pseudospin by a magnetic field could result in a sign change in the case of the presence of a Dzyaloshinskii-Moriya interaction (DMI) [30], which is not expected in NiO. The lack of a DMI also precludes the domain wall from acting as a polarizer or a retarder as described by Lan et al. [31], which thus cannot be a reason for a negative SSE sign in NiO.

It is not fully understood why this bulk NiO shows both a sign change as the
following increase in signal strength with distance. The difference between FMs and easy-axis AFMs, however, is that in AFMs there are domains present at the relevant field strengths. The partial reflection and/or absorption of magnons at domain walls could give a similar upturn of $\mu_m$ as the boundary conditions at the FM/PM interface. While $S_S$ is altered by the magnetic field strength, this could affect $\mu_m$ such that the region where the sign change occurs is stretched. The possibility exists that the maximum reached after the sign change is then shifted towards further distances such that the signal is increasing with distance in the investigated length scale. In addition to providing an explanation for the increase in $\mu_m$, the magnon accumulation at domain walls can be responsible for the large fluctuations between data points which increase with increasing distance due to movement of domain walls creating local variation in $\mu_m$.

Figure 6.4(c) shows the magnetic field strength dependence, fitted with the same models as are used for the local data, assuming that the $S_S$ is dominant for the change in signal strength. Also for the nonlocal signals the approach described in Ref. [15] using the dispersion from Ref. [16] best resembles the near-linear field dependence without offset. The SSE signal is driven by Joule heating and therefore expected to have a quadratic dependence on the current sent through the injector. However, after an initial increase in the current dependence, a saturation sets in around 75 $\mu$A as shown in Fig. 6.4(d). A strictly quadratic dependence is only expected when the parameters of the system do not change. Rising temperatures due to the current could lower the strongly temperature-dependent signal strength, leading to the saturation. The temperature increase could lower the spin Seebeck coefficient of the dependence of the magnon mode energies and their splitting by the magnetic field strength. Furthermore, an increased heat conductivity at the elevated temperature will create a lower temperature gradient with the same heating power. Both the field dependence and the absence of an electrically injected signal resembles the results reported for Cr$_2$O$_3$ by Yuan et al. [5]. However, they claim this spin transport is a result of spin superfluidity.

The SSE originates from the influence of a magnetic field on the population of the magnon modes, but these models might be influenced by the movement of domain walls. The lacking offset in method B could be explained by the multi-domain nature as the domains have opposite magnon spin polarization resulting in a smaller net SSE. Moreover, domain walls can interact with the spin current itself, leading to domain wall movement [32] and spin current reflection upon domain walls [33]. In thin films with multiple domain walls, the reflection damps the nonlocal signals, but domain optimization by tuning the growth direction or by magnetic training still leads to micrometer spin transport [34].
6.7 Conclusion

To conclude, we observed a spin Seebeck effect generated spin current in the bulk easy-plane antiferromagnet NiO. This was achieved as a result of an applied magnetic field without the need of exchange interactions to align the magnetic domains. The field dependence of the SSE amplitude at 5K was modeled by the energy splitting of magnon modes, creating an imbalance in the magnon spin population. The cubic anisotropy and magnetic dipole-dipole interactions have to be taken into account in order to recreate the near-linear SSE dependence on the field and the small offset. Furthermore, the SSE signal exhibits both a sign change and then an increase with increasing the injector distance that would not be possible without the introduction of additional boundary conditions in the bulk NiO, a role that may have been fulfilled by domain walls.

6.8 Outlook

Since the paper version of this chapter was submitted, further developments in the field took place. Ref. [30] [35] pointed out that there is controversy around whether magnons in easy-plane AFMs carry magnetic momentum in general. May this not be the case, then this has implications for some parts of the interpretation and the model in this chapter. Furthermore, although lacking DMI, the anisotropy in NiO could result in the rotation of the pseudo-spin. The precession frequency is given by

\[ h\Omega = h\omega_{an} - \mu_0 H_{DMI} m_{net}, \]

where \( \Omega \) is the pseudo spin precession frequency, \( \omega_{an} \) the normalized anisotropy frequency, \( m_{net} \) the net magnetization due to the DMI and \( H_{DMI} \) is the effective DMI field. As shown by Eq. (6.5), the magnetic field strength has little effect on this rotation because \( H_{DMI} \) is not a function of the magnetic field and \( m_{net} \) is small. Nevertheless, the pseudospin could be rotated as a function of the distance. Different domains could act differently on this rotation, having a different \( \omega_{an} \). This could be another explanation for the sign change observed in the distance dependence as well as the increase in error bars with increasing distance. Therefore, the main conclusion of this chapter holds.
Bibliography

[1] A. V. Chumak, V. I. Vasyuchka, A. A. Serga, and B. Hillebrands, “Magnon spintronics,” *Nature Physics* **11**(June), pp. 1505–1549, 2015.

[2] L. J. Cornelissen, J. Liu, R. A. Duine, J. Ben Youssef, and B. J. Van Wees, “Long-distance transport of magnon spin information in a magnetic insulator at room temperature,” *Nature Physics* **11**, p. 1022, 2015.

[3] R. Lebrun, A. Ross, S. A. Bender, A. Qaiumzadeh, L. Baldrati, J. Cramer, A. Brataas, R. A. Duine, and M. Kläui, “Tunable long-distance spin transport in a crystalline antiferromagnetic iron oxide,” *Nature* **561**, p. 222, 2018.

[4] W. Xing, L. Qiu, X. Wang, Y. Yao, Y. Ma, R. Cai, S. Jia, X. C. Xie, and W. Han, “Magnon Transport in Quasi-Two-Dimensional van der Waals Antiferromagnets,” *Physical Review X* 9(1), p. 11026, 2019.

[5] W. Yuan, Q. Zhu, T. Su, Y. Yao, W. Xing, Y. Chen, Y. Ma, X. Lin, J. Shi, R. Shindou, X. C. Xie, and W. Han, “Experimental signatures of spin superfluid ground state in canted antiferromagnet Cr$_2$O$_3$ via nonlocal spin transport,” *Science Advances* **4**, p. 1098, 2018.

[6] A. Qaiumzadeh, H. Skarsvåg, C. Holmqvist, and A. Brataas, “Spin Superfluidity in Biaxial Antiferromagnetic Insulators,” *Physical Review Letters* **118**(13), p. 137201, 2017.

[7] L. J. Cornelissen and B. J. Van Wees, “Magnetic field dependence of the magnon spin diffusion length in the magnetic insulator yttrium iron garnet,” *Physical Review B* 93(R), p. 020403, 2016.

[8] K. Shen, “Pure spin current in antiferromagnetic insulators,” *Physical Review B* **100**(9), p. 94423, 2019.

[9] B. Flebus, “Chemical potential of an antiferromagnetic magnon gas,” *Physical Review B* **100**, p. 064410, 2019.

[10] S. M. Rezende, R. L. Rodríguez-Suárez, and A. Azevedo, “Theory of the spin Seebeck effect in antiferromagnets,” *Physical Review B* **93**, p. 014425, 2016.

[11] S. M. Wu, W. Zhang, A. Kc, P. Borisov, J. E. Pearson, J. S. Jiang, D. Letterman, A. Hoffmann, and A. Bhattacharya, “Antiferromagnetic Spin Seebeck Effect,” *Physical Review Letters* **116**, p. 097204, 2016.

[12] R. E. Troncoso, S. A. Bender, A. Brataas, and R. A. Duine, “Spin transport in thick insulating antiferromagnetic films,” *Physical Review B* **101**, p. 054404, 2020.

[13] J. Shan, L. J. Cornelissen, N. Vlietstra, J. B. Youssef, T. Kuschel, R. A. Duine, and B. J. Van Wees, “Influence of yttrium iron garnet thickness and heater opacity on the nonlocal transport of electrically and thermally excited magnons,” *Physical Review B* **94**, p. 174437, 2016.
[14] J. Holanda, D. S. Maior, O. Alves Santos, L. H. Vilela-Leão, J. B. Mendes, A. Azevedo, R. L. Rodríguez-Suárez, and S. M. Rezende, “Spin Seebeck effect in the antiferromagnet nickel oxide at room temperature,” Applied Physics Letters 111, p. 172405, 2017.

[15] S. M. Rezende, A. Azevedo, and R. L. Rodríguez-Suárez, “Introduction to antiferromagnetic magnons,” Journal of Applied Physics 126, p. 151101, 2019.

[16] J. Milano and M. Grimsditch, “Magnetic field effects on the NiO magnon spectra,” Physical Review B 81, p. 094415, 2010.

[17] M. T. Hutchings and E. J. Samuelsen, “Measurement of spin-wave dispersion in NiO by inelastic neutron scattering and its relation to magnetic properties,” Physical Review B 6(9), p. 3447, 1972.

[18] T. Kampfrath, A. Sell, G. Klatt, A. Pashkin, S. Mährlein, T. Dekorsy, M. Wolf, M. Fiebig, A. Leitenstorfer, and R. Huber, “Coherent terahertz control of antiferromagnetic spin waves,” Nature Photonics 5(1), p. 31, 2011.

[19] M. Grimsditch and L. McNeil, “Unexpected behavior of the antiferromagnetic mode of NiO,” Physical Review B 58(21), p. 14462, 1998.

[20] P. R. T. Ribeiro, F. L. A. Machado, M. Gamino, A. Azevedo, and S. M. Rezende, “Spin Seebeck effect in antiferromagnet nickel oxide in wide ranges of temperature and magnetic field,” Physical Review B 99, p. 094432, 2019.

[21] N. Vlietstra, J. Shan, B. J. van Wees, M. Isasa, F. Casanova, and J. B. Youssef, “Simultaneous detection of the spin-Hall magnetoresistance and the spin-Seebeck effect in platinum and tantalum on yttrium iron garnet,” Physical Review B 90, p. 174436, 2014.

[22] G. R. Hoogeboom, A. Aqeel, T. Kuschel, T. T. M. Palstra, and B. J. van Wees, “Negative spin Hall magnetoresistance of Pt on the bulk easy-plane antiferromagnet NiO,” Applied Physics Letters 111, p. 052409, 2017.

[23] J. Milano, L. B. Steren, and M. Grimsditch, “Effect of dipolar interaction on the antiferromagnetic resonance spectra of NiO,” Physical Review Letters 93, p. 077601, 2004.

[24] W. L. Roth, “Neutron and optical studies of domains in NiO,” Journal of Applied Physics 31(11), p. 2000, 1960.

[25] J. Fischer, O. Gomonay, R. Schlitz, K. Ganzhorn, N. Vlietstra, M. Althammer, H. Huebl, M. Opel, R. Gross, S. T. B. Goennenwein, and S. Geprägs, “Spin Hall magnetoresistance in antiferromagnet / heavy-metal heterostructures,” Physical Review B 97, p. 014417, 2018.

[26] I. Sänger, V. V. Pavlov, M. Bayer, and M. Fiebig, “Distribution of antiferromagnetic spin and twin domains in NiO,” Physical Review B 74, p. 144401, 2006.

[27] S. M. Rezende, A. Azevedo, and R. L. Rodríguez-Suárez, “Magnon diffusion theory for the spin Seebeck effect in ferromagnetic and antiferromagnetic insulators,” Journal of Physics D: Applied Physics 51, p. 174004, 2018.
[28] J. Holanda, O. A. Santos, R. O. Cunha, J. B. S. Mendes, A. Azevedo, and S. M. Rezende, “Longitudinal spin Seebeck effect in permalloy separated from the anomalous Nernst effect: Theory and experiment,” Physical Review B 95, p. 214421, 2017.

[29] N. Vlietstra, J. Shan, V. Castel, J. Ben Youssef, G. E. W. Bauer, and B. J. van Wees, “Exchange magnetic field torques in YIG/Pt bilayers observed by the spin-Hall magnetoresistance,” Applied Physics Letters 103, p. 032401, 2013.

[30] T. Wimmer, A. Kamra, J. Gückelhorn, M. Opel, S. Geprägs, R. Gross, H. Huebl, and M. Althammer, “Coherent control of magnonic spin transport in an antiferromagnetic insulator,” Physical Review Letters 125, p. 247204, 2020.

[31] J. Lan, W. Yu, and J. Xiao, “Antiferromagnetic domain wall as spin wave polarizer and retarder,” Nature Communications 8, p. 178, 2017.

[32] R. Wieser, E. Y. Vedmedenko, and R. Wiesendanger, “Indirect control of antiferromagnetic domain walls with spin current,” Physical Review Letters 106, p. 067204, 2011.

[33] E. G. Tveten, A. Qaiumzadeh, O. A. Tretiakov, and A. Brataas, “Staggered dynamics in antiferromagnets by collective coordinates,” Physical Review Letters 110, p. 127208, 2013.

[34] A. Ross, R. Lebrun, O. Gomonay, D. A. Grave, A. Kay, L. Baldrati, S. Becker, A. Qaiumzadeh, C. Ulloa, G. Jakob, F. Kronast, J. Sinova, R. Duine, A. Brataas, A. Rothschild, and M. Kläui, “Propagation Length of Antiferromagnetic Magnons Governed by Domain Configurations,” Nano Letters 20, p. 306, 2020.

[35] A. Kamra, T. Wimmer, H. Huebl, and M. Althammer, “Antiferromagnetic magnon pseudospin: Dynamics and diffusive transport,” Physical Review B 102(17), p. 174445, 2020.