Spin conductance in extended thin films of YIG driven from thermal to subthermal magnons regime by large spin-orbit torque

N. Thiery, A. Draveny, V. V. Naletov, L. Vila, J.P. Attané, C. Beigné, G. de Loubens, M. Viret, N. Beaulieu, J. Ben Youssef, V. E. Demidov, S. O. Demokritov, A. N. Slavin, V. S. Tiberkevich, A. Anane, P. Bortolotti, V. Cros, and O. Klein

1 SPINTEC, CEA-Grenoble, CNRS and Université Grenoble Alpes, 38054 Grenoble, France
2 Institute of Physics, Kazan Federal University, Kazan 420008, Russian Federation
3 SPEC, CEA-Saclay, CNRS, Université Paris-Saclay, 91191 Gif-sur-Yvette, France
4 LabSTICC, CNRS, Université de Bretagne Occidentale, 29238 Brest, France
5 Department of Physics, University of Muenster, 48149 Muenster, Germany
6 Institute of Metal Physics, Ural Division of RAS, Yekaterinburg 620041, Russian Federation
7 Department of Physics, Oakland University, Michigan 48309, USA
8 Unité Mixte de Physique CNRS, Thales, Université Paris-Saclay, 91767 Palaiseau, France

(Dated: September 19, 2017)

We report a study on spin conductance through ultra-thin extended films of epitaxial Yttrium Iron Garnet (YIG), where spin transport is provided by propagating spin waves, that are generated and detected by direct and inverse spin Hall effects in two Pt wires deposited on top. While at low current the spin conductance is dominated by transport of exchange magnons, at high current, the spin conductance is dominated by magnetostatic magnons, which are low-damping non-equilibrium magnons thermalized near the spectral bottom by magnon-magnon interaction, with consequent a sensitivity to the applied magnetic field. This picture is supported by microfocus Brillouin Light Scattering spectroscopy.

The recent demonstrations that spin orbit torques (SOT) allow one to generate and detect pure spin currents [1–8] has triggered a renewed effort to study magnons transport in extended magnetic films. This topic is currently recognized as one of the important emerging research direction in modern magnetism [9]. This is because, in contrast to spin transfer process in confined geometries (e.g. nano-pillars or nano-contacts) where usually the uniform magnon mode dominate the dynamics, very little is known about spin transfer in extended geometries, which have continuous spin-wave spectra containing many modes which can take part in the magnon-magnon interactions. A large effort has concentrated so far on yttrium iron garnet (YIG), a magnetic insulator, which is famous for having the lowest known magnetic damping parameter [10]. From a purely fundamental point of view, the studies of magnon transport in YIG by means of the direct and inverse spin Hall effects (ISHE) [2, 11–19] are very interesting, as they provide new means to alter efficiently the energy distribution of magnons and, potentially, even to trigger condensation [20].

Contrary to the case of magnons excited coherently, e.g. by means of a ferromagnetic resonance or parametric pumping, when the frequencies of the excited magnons are fully determined by the frequencies of the external signals, the excitation of magnons by means of spin transfer process lacks frequency selectivity [21], and, therefore, can lead to their excitation in a broad frequency range. This poses a challenge for the identification of the nature of magnons modes excited by SOT. It has been already shown in [22], that it is convenient and useful to introduce the concepts of subthermal (having energy close to the bottom of the spin wave spectrum) and thermal (having energy close to $k_B T$) magnons. On one hand, it has been well established [23, 24] that subthermal magnons can be very efficiently thermalized near the spectral bottom (region of so-called magnetostatic waves) by the intensive magnon-magnon interaction, whose decay rate between quasi-degenerate modes increases with power, to reach a quasi-equilibrium state by a non-zero chemical potential [23–25] and an effective temperature [26]. On the other hand, it has been shown, both experimentally and the-
oretically [26], that the groups of subthermal and thermal magnon are effectively decoupled from each other, as under the intensive parametric pumping one can reach a state where the effective temperature of subthermal magnons exceeds the real temperature characterizing the thermal magnons by a factor of 100.

Under spin transfer process, whose efficiency is known to increase with decreasing magnon frequency, in confined geometries with localized spin-current injection (i.e. when there are no quasi-degenerate modes), it has been shown, that one can reach current induced coherent GHz-frequency magnon dynamics in YIG [16, 17, 27]). In extended thin-films, the recently discovered non-local magnon transport [28–32] suggests that the magnon transport properties of YIG films subjected to small SOT are dominated by thermal magnons, whose number overwhelmingly exceeds the number of other modes at any non-zero temperature. The interesting challenge is to elucidate what will happen to this spectrum (in particular the interplay between the thermal and subthermal part [33]) when one applies large SOT to a magnon continuum.

We propose herein to measure the room temperature spin conductance of YIG films when the driving current is varied in a wide range magnitudes [34, 35] creating, first, a quasi-equilibrium transport regime, and, then, driving the system to a strongly out-of-equilibrium state. To reach this goal the spin current injected in the YIG by SOT shall be increased by more than one order of magnitude compared to previous works, while simultaneously reducing the film thickness by also an order of magnitude, using ultra-thin films of YIG grown by liquid phase epitaxy (LPE) [14, 36]. A series of lateral devices have been patterned on a 18 nm thick YIG films. Ferromagnetic resonance (FMR) characterization of the bare film are summarized in Table 1. On these films, we have deposited Pt wires, 10 nm thick, 300 nm wide, and 20 μm long. The measured resistance of the Pt wire at room temperature is $R_0 = 1.3 \, \text{kΩ}$. The lateral device geometry is shown in FIG.1a. One monitors the voltage $V$ along one wire as a current $I$ flows through a second wire separated by a gap of 1.2 μm. Here the Pt wires are connected by 50 nm thick Al electrodes colored in yellow. Since large amount of electrical current needs to flow in the Pt, a pulse method is used to reduce significantly Joule heating. In the following the current is injected during 10 ms pulses series enclosed in a 10% duty cycle. Temperature sensing is provided by the change of relative resistance of the Pt wire during the pulse. In FIG.1b, we have plotted $\kappa_\text{Pt}(R_I - R_0)/R_0$ as a function of the current $I$, where the coefficient $\kappa_\text{Pt} = 254 \, \text{K}$ is specific to Pt. We observe that the pulse method allows to keep the absolute temperature of our YIG below 340 K [37] at the maximum current amplitude of 2.5 mA. Avoiding excessive heating of the YIG is crucial because, in a joint review paper [38], it is shown that epitaxial YIG films grown by LPE behave as a large gap semiconductor, with an electrical resistivity that decreases exponentially with increasing temperature following an activated behavior. As shown in Ref. [38], at 340 K, however, the electrical resistivity of YIG remains larger than $10^6 \, \text{Ω}\cdot\text{cm}$ and thus the YIG can still be considered a good insulator ($R > 30 \, \text{GΩ}$) over the current range explored herein.

The lateral device is biased by an in-plane magnetic field, $H_0$ set at a variable polar angle $\theta$, with respect to the perpendicular of the Pt wires. FIG.2a–d displays the results when $I = 1.5 \, \text{mA}$ and $H_0 = 2 \, \text{kOe}$ [39]. For each value of $\theta$, 4 measurements $V_\pm^\pm$ are performed corresponding to the 4 combinations of the polarities of $\pm H_0$ and $\mp I$ (the polarity convention is defined in FIG.1a).

| YIG  | $t_{\text{YIG}}$ (nm) | $4\pi M_s$ (G) | $\alpha_{\text{YIG}}$ | $\Delta H_0$ (Oe) |
|------|----------------------|---------------|------------------|-----------------|
| 18   | $1.6 \times 10^4$    | $4.4 \times 10^{-3}$ | 3.7 | 
| Pt   | $t_{\text{Pt}}$ (nm) | $\rho$ (μΩ cm) | $\alpha_{\text{YIG}/\text{Pt}}$ | $g_{\text{Pt}}$ (m$^{-2}$) |
| 10   | 19.5                 | $2.4 \times 10^{-3}$ | $3 \times 10^{18}$ |

| TABLE I. Summary of the physical properties of the materials used in this study. |

---

FIG. 2. (Color online) Angular dependence of the non-local voltages $V_\pm^\pm$ measured while inverting the polarity of the applied field $H_0 = \pm 2 \, \text{kOe}$ (red/blue) respectively for a) negative and b) positive current pulses $I = \mp 1.5 \, \text{mA}$. The measured signal can be decomposed c) and d) in three components: $\Sigma$ (green): the signal sum, $\Delta$ (orange): the signal difference and $V_j$: the offset; respectively even/odd, odd/even in field/current, and an independent contribution (dashed). Panel e) shows the current dependence of the amplitude $\Sigma$ and $\Delta$.
FIG. 2a and 2b show the raw data obtained respectively for negative and positive current pulses. Clearly the non-local voltage oscillates around an offset, \(V_\parallel\), defined as the voltage measured at \(\theta = \pm 90^\circ\). This offset is independent of the current polarity and its amplitude scales with the temperature elevation of Pt produced by Joule heating (see FIG.2c in [38]). We ascribe it to thermoelectric effects produced by a small temperature difference at the two Pt|Al contacts of the detector circuit [40]. By contrast, the anisotropic part of the voltage is ascribed to magnons transport.

To gain more insight, the data are sorted according to their symmetry with respect to the \((H_\|, I)\)-polarity. This is done in FIG.2e and 2d by constructing the signal sum \(\Sigma_\| = (V^+ + V_-)/2 - V_\parallel\) (green tone) even in field and the signal difference \(\Delta_\| = (V^+ - V_-)/2\) (orange tone) odd in field. This separation is exposed in their angular dependences, which follow two different behaviors, one in \(\cos^2\theta\), the other one in \(\cos\theta\), respectively. The solid lines in FIG.2e and in 2d are fit by these two functions. Comparing the behavior, we observe that the signal \(\Sigma\) is odd in current, while the signal \(\Delta\) is even in current. As noticed in ref [28], these symmetries of \(\Sigma\) and \(\Delta\) are the hallmark of respectively SOT [17] and spin Seebeck effects [41, 42]. Hereafter, we shall use the fit of the whole angular dependence as a mean to extract precisely the amplitude of \(\Sigma\) and \(\Delta\) at \(\theta = 0\).

FIG.2e shows their evolution as a function of current. One observes that the correspondence between the symmetries of \(\Sigma\) and of \(\Delta\) with respect to the polarities of \(H_\|\) and \(I\) is respected (within our measurement accuracy) on the whole current range. While \(\Delta\) approximately follows the parabolic increase of the Pt temperature (cf. FIG.1a), as expected for thermal effects, the interesting novel feature is the fact that \(\Sigma\) deviates from a purely linear transport behavior at large \(I\). It is important also to notice that, when the high/low binding posts of the current source and voltmeter are biased in the same orientation (cf. FIG.1a), the sign of \((\Sigma \cdot I)\) < 0. This is a signature that the observed non-local voltage is produced by ISHE and not by leakage electrical currents inside the YIG [38], although these effects are only expected to occur at much higher temperatures (> 370 K[43]). While in both scenario the induced electrical current flows in the same direction in the two parallel Pt wires, for ISHE, the YIG acts as a source and the potential increases along the current direction [44], in contrast, for Ohmic loss, the YIG acts as a load and the potential drops along the current direction (i.e. \((\Sigma \cdot I) > 0\) cf. FIG2b in [38]). We should though add that selecting the component of the non-local voltage that is \(\theta\)-dependent is another effective mean to eliminate Ohmic contribution, since the later are independent of the in-plane orientation of \(H_\|\). We have repeated this measurement on other devices either on the same film or on different LPE YIG films of similar thickness. On all the devices, we observe an up-turn of \(\Sigma\) at the same current density. While the sign of \((\Sigma \cdot I)\) is always negative, this is not the case for the sign for \((\Delta \cdot I)\), which depends on the film quality (the same \((\Delta \cdot I)\) sign is observed for all devices on the same film but could change depending on the film quality). Further progress on the later issue requires a better understanding of the different phenomena contributing to thermal effects \((\Delta\text{-signal})\) and the means to separate them. In the following, we shall concentrate exclusively on the non-linear behavior of \(\Sigma\) which measures the number of magnons created by SOT relatively to the number of magnons annihilated by SOT while being immune to effects caused by Joule heating.

Using open dots, we have plot in FIG.3a both the variation of \(|\Sigma^-|\) and \(|\Sigma^+|\) as a function of the current intensity. Both data set show clearly the emergence of a new spin transport channel at large current densities, as evidence by the deviation from a purely linear conduction regime. Since both quantities \(|\Sigma^\pm|\) follow the same behavior on the whole current range, for the sake of simplicity we shall call simply \(\Sigma\) (dark green) their averaged. At low current, the SOT signal follows first a linear behavior which has been shown to be dominated by thermal magnons transport [28] [45]. We shall define \(\Sigma^{(t)}\) the additional conduction contribution, i.e. the deviation from the extrapolated linear behavior.

Quite remarkably the enhancement of the conductance due to \(\Sigma^{(s)}\) occurs very gradually. We emphasize that such a low rise is very different from the sudden surge of coherent magnons observed at the critical threshold in confined geometries [46, 47]. Fitting a straight line through the low current regime, \(\Sigma_{I\epsilon[0,0.9]\text{mA}}\), and the high current deviation, \(\Sigma_{I\epsilon[1.6,2.3]\text{mA}}\), the intersection provides...
an estimation of the onset current of this new conduction channel, \( I_c \approx \pm 1.5 \, \text{mA} \), which corresponds to a current density \( J_c \approx 5 \times 10^{11} \, \text{A/m}^2 \). This value is very close to the threshold current for damping compensation of coherent modes observed at the same applied field \( (H_0 = 2 \, \text{kOe}) \) in micron-sized disks [17] and stripes [48] with similar characteristics.

More insight about the nature of the magnons excited above \( I_c \) can be obtained by studying the field dependence of \( \Sigma \) [49]. The results are shown in FIG 3b and 3c for two values of the current \( I = 0.4 \) and \( 2.5 \, \text{mA} \), respectively below and above \( I_c \). While in the field range explored, the signal is almost independent of \( H_0 \) when \( I < I_c \), it becomes strongly field dependent when \( I > I_c \). This different behaviors are consistent with assigning the spin transport to thermal magnons below \( I_c \) and mainly to subthermal magnons above \( I_c \). In the former case, the magnons’ energy is of the order of the exchange energy, which is much larger than the Zeeman energy, while in the latter case, because of their long wavelength, their energy is of the order of the magnetostatic energy. In consequence, \( \Sigma \) is expected to increase with decreasing field at fixed \( I \), because of the associated decrease of \( I_c \). The behavior scales well with the reduced quantity \( I/I_c \). This is shown by the solid line in FIG 3b, which displays the expected field dependence of \( 1/I_c(H_0) \) [17]) where \( I_c \propto (\omega_H + \omega_M/2)(\alpha + \gamma \Delta H_0/(2\omega_K)) \), where \( \omega_H = \gamma H_0 \) and \( \omega_M = \gamma M \), \( \gamma \) being the gyromagnetic ratio, and \( \omega_K = \sqrt{\omega_H(\omega_H + \omega_M)} \) is the Kittel’s law. We have used here the amount of inhomogeneous broadening, \( \Delta H_0 = 1.5G \) (probably position dependent), as an adjustable parameter, while the value of the other parameters are those extracted from Table 1.

The above interpretation has been checked by performing microfocus Brillouin light scattering (\( \mu \)-BLS) in the sub-thermal energy range. For this measurement, we have used a second series [50] of non-local devices, where the Pt thickness has been reduced to 7 nm (thus comparison of the results between the 2 series should be done by juxtaposing data obtained with indentical current densities, cf. upper scale). FIG 4a and 4b show on a logarithmic scale the spectral distribution of the BLS intensity, \( J \), a) underneath the injector and b) underneath the detector, which are here separated by \( d = 0.7 \mu \text{m} \). The distribution is measured at \( I = \pm 2 \, \text{mA} \) (i.e. \( 9.5 \times 10^{11} \, \text{A/m}^2 \)) while the field is set to \( H_0 = \pm 2 \, \text{kOe} \). In both cases, an enhancement of the subthermal magnons population is observed when \( (I \cdot H_0) < 0 \) (blue), which corresponds to the configuration where the SOT compensates the damping (cf convention in FIG1a). The measurement for the opposite case \( (I \cdot H_0) > 0 \) (red) provides a reference about the out-of-equilibrium state produced by Joule heating. The maximum intensity of the red curve indicates the resonance frequency of the Kittel mode, \( \omega_K/2\pi \) at the corresponding temperature. This is because the \( \mu \)-BLS response function is centered around the long wavelength magnons. Indeed, the detected signal decreases once the magnon wavelength is smaller than the spot size (approximately \( 0.4 \mu \text{m} \): diffraction limited).

In order to isolate the contribution produced by SOT, we subtract the spectral contribution measured at \( +I \) to the one measured at \( -I \) (grey shaded area). This allows us to cancel out the spectral deformation produced by Joule heating but, as for the \( \Sigma \)-signal, this only measures the enhancement of the magnons created by SOT relative to the magnons annihilated by SOT. One can clearly see on the shaded data that SOT enhances the magnons population in a spectral window between the Kittel frequency and the bottom of the magnon manifold. Next, we have plotted in FIG4c how the spectral integration of this differential signal \( J_\pm = \int J_\Sigma \, d\omega \) varies as a function of the current amplitude underneath the injector. One observes a regime of linear rise at small current, followed by a growth above \( J_c \approx 5 \times 10^{11} \, \text{A/m}^2 \) in a similar fashion as the one reported in FIG3a. The \( \mu \)-BLS experiment thus provides a direct evidence that an additional spin conduction channel has indeed emerged in the GHz frequency range (subthermal) at large current when SOT is in the range to compensate the damping. It also shows that the magnons newly created are spread at the bottom of the magnon manifold.

In summary, we have shown that while at low values of the spin current the main contribution to the spin conductance comes from thermal magnons, the subthermal magnons mainly determine the magnon transport at high values of the spin current, comparable to the critical magnitude at which damping compensation of coherent magnons takes place. We believe that our current findings are not only important from the fundamental viewpoint of view, but might be also useful for future
applications. While transport of thermal magnons are difficult to control due to their relatively high energies, the subthermal magnons could be efficiently controlled by variation of relatively weak magnetic fields.

This research was supported in part by the CEA program NanoScience (project MAFETY), by the priority program SPP1538 Spin Caloric Transport (SpinCat) of the DFG and by the program Megagrant 14.Z50.31.0025 of the Russian ministry of Education and Science. The work at Oakland University was supported by the Grants Nos. EFMA-1641989 and ECCS-1708982 from the NSF of the USA, by the CND, NRI and by DARPA. VVN acknowledges fellowship from the emergence strategic program of UGA, and Russian competitive growth program. We thank G. Zhand, T. van Pham, A. Brenac for their help in the fabrication of the lateral devices.

* Corresponding author: oklein@cea.fr

[1] S. O. Valenzuela and M. Tinkham, Nature 442, 176 (2006).
[2] Y. Kajiwara, K. Harii, S. Takahashi, J. Ohe, K. Uchida, M. Mizuguchi, H. Umezawa, H. Kawai, K. Aono, K. Takahashi, S. Maekawa, and E. Saitoh, Nature 404, 262 (2000).
[3] I. M. Miron, K. Garello, G. Gaudin, P.-J. Zermatten, S. Sangiao, J. M. De Teresa, L. Morellon, I. Lucas, M. C. Martinez-Velarte, and A. Fert, Nature Comm. 4, 2944 (2013).
[4] S. Sangiao, J. M. De Teresa, L. Morellon, I. Lucas, M. C. Martinez-Velarte, and M. Viret, Appl. Phys. Lett. 106, 172403 (2015).
[5] A. R. Mellnik, J. S. Lee, A. Richardella, J. L. Grab, P. J. Mintun, M. H. Fischer, A. Vaezi, A. Manchon, E.-A. Kim, N. Samarth, and D. C. Ralph, Nature 511, 449 (2014).
[6] E. Lesne, Y. Fu, S. Oyarzun, J. C. Rojas-Sánchez, D. C. Vaz, H. Nagamura, G. Sicoli, J.-P. Attan, M. Jamet, E. Jacquet, J.-M. George, A. Barthlmy, H. Jaffrs, A. Fert, M. Bébes, and L. Vila, Nature Mater. 15, 1261 (2016).
[7] J.-Y. Chauleau, M. Boselli, S. Gariglio, R. Weil, G. de Loubens, J.-M. Triscone, and M. Viret, EPL 116, 17006 (2016).
[8] D. Sander, S. O. Valenzuela, D. Makarov, C. H. Marrows, E. E. Fullerton, P. Fischer, J. McCord, P. Vavassori, S. Mangin, P. Pirro, and et al., Journal of Physics D: Applied Physics 50, 363001 (2017).
[9] E. G. Spencer, R. C. LeCraw, and A. M. Clogston, Phys. Rev. Lett. 3, 32 (1959).
[10] Z. Wang, Y. Sun, M. Wu, V. Tiberkevich, and A. Slavin, Phys. Rev. Lett. 107, 146601 (2011).
[11] E. Padrón-Hernández, A. Azevedo, and S. M. Rezende, Appl. Phys. Lett. 99, 192511 (2011).
[12] A. V. Chumak, A. A. Serga, M. B. Jungfleisch, R. Neb, D. A. Bozhko, V. S. Tiberkevich, and B. Hillebrands, Appl. Phys. Lett. 100, 082405 (2012).
[13] C. Hahn, G. de Loubens, O. Klein, M. Viret, V. V. Natalov, and J. Ben Youssef, Phys. Rev. B 87, 174417 (2013).
[14] O. d’Allivy Kelly, A. Anane, R. Bernard, J. Ben Youssef, C. Hahn, A. H. Molpeceeres, C. Carretero, E. Jacquet, C. Deranlot, P. Bortolotti, R. Lebourgeois, J.-C. Mage, G. de Loubens, O. Klein, V. Cros, and A. Fert, Appl. Phys. Lett. 103, 082408 (2013).
[15] A. Hamadeh, O. d’Allivy Kelly, C. Hahn, H. Meley, R. Bernard, A. H. Molpeceeres, V. V. Naletov, M. Viret, A. Anane, V. Cros, S. O. Demokritov, J. L. Prieto, M. Muñoz, G. de Loubens, and O. Klein, Phys. Rev. Lett. 113, 197205 (2014).
[16] M. Collet, X. de Milly, O. d’Allivy Kelly, V. Naletov, R. Bernard, P. Bortolotti, J. Ben Youssef, V. Demidov, S. Demokritov, J. Prieto, M. Muñoz, V. Cros, A. Anane, G. de Loubens, and O. Klein, Nature Commun. 7, 10377 (2016).
[17] V. Lauer, D. A. Bozhko, T. Bricher, P. Pirro, V. I. Vasyuchka, A. A. Serga, M. B. Jungfleisch, M. Agrawal, Y. V. Koblijanskyj, G. A. Melkov, C. Dubs, B. Hillebrands, and A. V. Chumak, Appl. Phys. Lett. 108, 012402 (2016).
[18] D. Wesenberg, T. Liu, D. Balzar, M. Wu, and B. L. Zink, Nat Phys advance online publication, (2017).
[19] S. A. Bender, R. A. Duine, and Y. Tserkovnyak, Phys. Rev. Lett. 108, 246601 (2012).
[20] V. E. Demidov, S. Urazhdin, E. R. J. Edwards, M. D. Stiles, R. D. McMichael, and S. O. Demokritov, Phys. Rev. Lett. 107, 107204 (2011).
[21] V. E. Demidov, S. Urazhdin, E. R. J. Edwards, M. D. Stiles, R. D. McMichael, and S. O. Demokritov, Phys. Rev. Lett. 107, 107204 (2011).
[22] V. E. Demidov, O. Dzyapko, G. A. Melkov, A. A. Serga, B. Hillebrands, and A. N. Slavin, Nature 443, 430 (2006).
[23] V. E. Demidov, O. Dzyapko, G. A. Melkov, A. A. Serga, B. Hillebrands, and A. N. Slavin, Phys. Rev. Lett. 99, 037205 (2007).
[24] C. Du, T. van der Sar, T. X. Zhou, P. Upadhyaya, F. Casola, H. Zhang, M. C. Onbasli, C. A. Ross, R. L. Walsworth, Y. Tserkovnyak, and et al., Science 357, 195198 (2017).
[25] A. A. Serga, V. S. Tiberkevich, C. W. Sandweg, V. I. Vasyuchka, D. A. Bozhko, A. V. Chumak, T. Neumann, B. Obry, G. A. Melkov, A. N. Slavin, and B. Hillebrands, Nature Comm 5, 3452 (2014).
[26] V. E. Demidov, M. Evelt, V. Bessonov, S. O. Demokritov, J. L. Prieto, M. Muoz, J. Ben Youssef, V. V. Naletov, G. de Loubens, O. Klein, M. Collet, P. Bortolotti, V. Cros, and A. Anane, Sci. Rep. 6, 32781 (2016).
[27] L. J. Cornelissen, J. Liu, R. A. Duine, J. Ben Youssef, and B. J. van Wees, Nature Physics 11, 1022 (2015).
[28] S. T. B. Goennenwein, R. Schlitz, M. Perspinitier, K. Ganzhorn, M. Althammer, R. Gross, and H. Huebl, Appl. Phys. Lett. 107, 172405 (2015).
[29] J. Li, Y. Xu, M. Aldesary, C. Tang, Z. Lin, S. Zhang, L. Wang, and J. Shi, Nature Commun. 7, 10858 (2016).
[30] H. Wu, C. H. Wan, X. Zhang, Z. H. Yuan, Q. T. Zhang, J. Y. Qin, H. X. Wei, X. F. Han, and S. Zhang, Phys. Rev. B 93, 060403 (2016).
[31] L. J. Cornelissen, K. J. H. Peters, G. E. W. Bauer, R. A. Duine, and B. J. van Wees, Phys. Rev. B 94, 014412 (2016).
[32] S. Lendinez, J. Hang, S. Vélez, J. M. Hernández, D. Backes, A. D. Kent, and F. Mačić, Phys. Rev. Applied 7, 054027 (2017).
YIG spontaneous magnetization decreases by about 4 G/°C.

This corresponds to injecting in the Pt wire current densities > 1.0 × 10^{12} A-m^2.

This effect is independent of the sign of the spin Hall angle.

The non-local linear resistance between the two Pt wires is \( \Sigma^{(t)}/I = 0.019 \) mΩ.

This second series has been used to investigate the electrical properties of YIG thin films at high temperature.