Spin Hall noise

Akashdeep Kamra,1,2 Friedrich P. Witek,1 Sibylle Meyer,1,4 Hans Huebl,1,3 Stephan Geprägs,1 Rudolf Gross,1,3,4 Gerrit E. W. Bauer,1,5,6 and Sebastian T. B. Goennenwein1,3

1Walther-Meißner-Institut, Bayerische Akademie der Wissenschaften, 85748 Garching, Germany
2Kavli Institute of NanoScience, Delft University of Technology, 2628 CJ Delft, The Netherlands
3Nanosystems Initiative Munich (NIM), 80799 Munich, Germany
4Physik-Department, Technische Universität München, 85748 Garching, Germany
5Institute for Materials Research, Tohoku University, Sendai 980-8577, Japan
6WPI Advanced Institute for Materials Research, Tohoku University, Sendai 980-8577, Japan

(Rceived 30 September 2014; revised manuscript received 20 November 2014; published 8 December 2014)

We measure the low-frequency thermal fluctuations of pure spin current in a platinum film deposited on yttrium iron garnet via the inverse spin Hall effect (ISHE)-mediated voltage noise as a function of the angle $\alpha$ between the magnetization and the transport direction. The results are consistent with the fluctuation-dissipation theorem in terms of the recently discovered spin Hall magnetoresistance (SMR). We present a microscopic description of the $\alpha$ dependence of the voltage noise in terms of spin-current fluctuations and ISHE.

DOI: 10.1103/PhysRevB.90.214419 PACS number(s): 72.25.Mk, 72.70.+m

I. INTRODUCTION

The quote “The noise is the signal” by Rolf Landauer [1] emphasizes the usefulness of noise spectroscopy in gaining deeper insight into physical phenomena ranging from astronomical [2] to mesoscopic [3–5] scales. The voltage fluctuations across a resistor in thermal equilibrium, known as the Johnson-Nyquist (JN) noise [6,7], is attributed to the charge current fluctuations due to the random thermal motion of the charge carriers in electrical conductors. It is much less appreciated that spin-current fluctuations exist in all metals since they do not interfere with most electronic processes. However, they become observable due to the spin-charge coupling in magnetic nanostructures [8,9]. The recently discovered spin Seebeck effect [10] is attributed to an imbalance of spin-current fluctuations [11,12] caused by a thermal gradient in a ferromagnet|normal metal bilayer system. Spin-dependent coherent transport could be detected in magnetic tunneling junctions (MTJs) via current shot-noise measurements [13,14]. However, a direct measurement of thermal spin-current noise has, to our knowledge, not been reported yet.

Here we report measurements of the voltage-noise power spectral density (PSD) and resistance across a platinum (Pt) thin film deposited on a yttrium iron garnet (YIG) layer as a function of the angle $\alpha$ between the applied magnetic field and the transport direction. These experiments are interpreted in terms of the thermal spin current noise in Pt modulated by the magnetization direction, which is transformed into charge noise by the inverse spin Hall effect (ISHE). The voltage PSD is found to obey the same angular dependence as the electric resistance, called spin Hall magnetoresistance (SMR) [15,16], consistent with the fluctuation-dissipation theorem (FDT) [17]. Since spin Hall effect (SHE) [18,19] is believed to be the dominant spin-charge coupling mechanism in heavy-metals films, we refer to our measurements as “spin Hall noise” (SHN).

The random thermal motion of the electrons in a normal metal (N) causes charge current, but because of their spin degree of freedom, also spin-current fluctuations. The ISHE converts spin-current into charge-current (or voltage) noise. In a ferromagnetic insulator (FI)|N heterostructure, the measured voltage noise $S_V = S'_V + S''_V$ is composed of spin-induced ($S'_V$) and charge ($S''_V$) noise. The FI modulates the conductor by selectively absorbing spin currents polarized normal to the magnetization direction, i.e., the spin transfer torque [20]. The implied dependence of the spin-induced noise power $S'_V$ on an applied magnetic field that controls the magnetization direction in the FI allows us to disentangle it from the charge noise $S''_V$ in the measured voltage noise $S_V$. A spin-charge coupling can, in principle, be achieved as well just at the FI|N interface by either the anomalous Hall effect in proximity-induced ferromagnetism in N [21] or a Rashba-type spin-orbit interaction [22]. However, there is evidence against a significant proximity effect at the YIG|Pt interface [23,24]. Furthermore, our basic result that the angular-dependent thermal noise is a direct measure of spin fluctuations is model-independent.

II. EXPERIMENTS

We first discuss the measurements of voltage noise and resistance of YIG|Pt bilayers. Samples were fabricated by depositing 60 nm of YIG ($Y_3Fe_5O_{12}$) on a 500-μm-thick, (111)-oriented gadolinium gallium garnet (Gd$_3$Ga$_5$O$_{12}$, GGG) substrate via pulsed laser deposition. A Pt film with thickness $t_N = 2.2$ nm was then grown in situ on top of the YIG film using electron beam evaporation. Subsequently, the sample was patterned into a Hall bar mesa structure (width $w = 80 \mu m$, length $l = 950 \mu m$) using optical lithography and argon ion beam milling. The detailed sample preparation is described in Ref. [25].

The voltage PSD was measured as sketched in Fig. 1(a). The voltage signal is fed into a Stanford Research fast Fourier transform spectrum analyzer (SR760) after amplification using a Stanford Research preamplifier (SR560). We refer to the square of the Fourier transform of a single finite duration time trace of the voltage signal as a “spectrum.” A “PSD

1Here the term “ferromagnets” includes ferrimagnets such as YIG.
FIG. 1. (Color online) (a) Schematic of the voltage power spectral density measurements. The sample (gray) is connected to a preamplifier and a FFT spectrum analyzer. The symbols + and − define the sign convention for the voltage measurements. The setup and the amplification stage are shielded by a metal box (red thick lines). The applied magnetic field (blue arrow) makes an angle \( \alpha \) with the voltage measurement direction. (b) A typical noise spectrum captured using the setup described in (a). The individual data points shown in Fig. 2(c) are averaged over the frequency window between 20 and 45 kHz. The dashed line depicts the white noise level expected from the fluctuation-dissipation theorem.

sweep” [as in Fig. 1(b)] is obtained by averaging 15,000 such spectra. A single average value of the white noise level is then obtained by averaging the PSD sweep data in the frequency range 20–45 kHz. The frequency window is so chosen in order to minimize the effects of the 1/f noise and external electromagnetic disturbances. The average of 19 such data points shown in Fig. 2(c) are averaged over the frequency window between 20 and 45 kHz. The dashed line depicts the white noise level expected from the fluctuation-dissipation theorem.

The measurement configuration is depicted in Fig. 2(a). A 60 mT magnetic field applied in the \( xz \) plane at an angle \( \alpha \) with the +z direction saturates the YIG magnetization along its direction. The voltage noise PSD \( S_{V,\text{long}} \) of the “longitudinal” voltage \( V_{\text{long}} \) [Fig. 2(a)] averaged over 19 \( \alpha \) sweeps is shown as white open squares in Fig. 2(c). We also carried out conventional SMR measurement [15] of the longitudinal resistance \( R_{\text{long}} \) along the Hall bar (\( zz \)) direction [Fig. 2(b)] as a function of \( \alpha \) for a charge current \( I_q = 40.5 \mu \text{A} \) along the Hall bar. \( R_{\text{long}} \), shown as red triangles in Fig. 2(c), exhibits the \( \cos^2 \alpha \)-dependence characteristic of the SMR effect [16]. We find that \( S_{V,\text{long}} \) and \( R_{\text{long}} \) are related by \( S_{V,\text{long}} = 4k_B T R_{\text{long}} \), with \( T = 291.5 \text{ K} \) (room temperature), as expected from the fluctuation-dissipation theorem. Since the \( \alpha \)-dependence of \( R_{\text{long}} \) is attributed to SHE-generated spin currents [16], the anisotropic PSD must be caused by the spin Hall noise.

III. THEORY

To substantiate this claim, in the following we present a statistical linear response theory for the \( \alpha \)-dependent noise that elucidates the role of the spin currents. We restrict the analysis to frequencies far below the ferromagnetic resonance (FMR) frequency \( f_0 \). We consider a bilayer of a normal metal (N) with spin Hall angle \( \theta_{\text{SH}} \) deposited on a ferromagnetic insulator (FI) with its equilibrium magnetization pointing along \( \hat{z} \) as shown in Fig. 3. The magnetization dynamics in the FI is described by the Landau-Lifshitz-Gilbert (LLG) equation:

\[
\dot{m} = -\gamma [m \times \mu_0 (H_{\text{eff}} + h_0)] + \alpha_0 (m \times \dot{m}),
\]

(1)

The noise floor of our setup (output with zero voltage input, i.e., short circuited amplifier input) \( 1.52 \times 10^{-17} \text{ V}^2/\text{Hz} \) is subtracted from all data points.
Gilbert damping constant, respectively. The effective magnetic
\[ \rho = \rho_{L} + \rho_{G} \]
coordinate system is depicted in red. The black arrows define our sign
\[ \mathbf{m} \]
component does not contribute to the ISHE signal [32], while
\[ \mathbf{J}_{x} = -\frac{\hbar \dot{\mathbf{G}}}{4\pi} m_{x} + M_{s} \mu_{0} h_{y} \]
with correlation function
\[ \langle J_{x}(q,t)J_{x}(q',t') \rangle = M_{s}^{2} \langle \mu_{x} h_{y}(q,t) \mu_{x} h_{y}(q',t') \rangle \]
Only the first term on the right-hand side of the equation above is appreciable [33] because the ac susceptibility, and therefore
\[ J_{x}(\mathbf{q},t)J_{x}(\mathbf{q}',t') = \frac{\hbar \dot{G}}{2\pi} k_{B} T \delta(t - t') \delta(\mathbf{q} - \mathbf{q}') \]
Using the above result and the Wiener-Khintchine theorem relating the one-sided PSD \( S_V(\omega) \) and the autocorrelation of a variable \( f(t) \), \( S_V(\omega) = 2 \int \langle f(t) f(0) \rangle e^{-i\omega t} dt \), PSD of the spin Hall noise reads
\[
S_V(\omega) = 2 \int \frac{b^2\tilde{g}_N}{2\pi} k_BT \omega \delta(t) e^{-i\omega t} dt,
\]
\[
= 4k_BT \rho_1 \frac{1}{w_{tN}} = 4k_BT R_1,
\]
where \( R_1 = \rho_1/\rho_{tN} \) and
\[
\rho_1/\rho = \left[ \theta_{SH} e^{2\lambda_{sd}} \tanh \left( \frac{t_N}{2\lambda_{sd}} \right) \right]^2 \frac{\tilde{g}_{e-p}}{4\pi h_N}. \tag{15}
\]
When the equilibrium magnetization direction makes an angle \( \alpha \) with the voltage measurement direction (\( \hat{z} \)), the right-hand side of Eq. (14) is simply multiplied by \( \cos^2 \alpha \) [33], because only the \( z \) projection (\( \cos \alpha \)) of the fluctuating ISHE current [Eq. (10)] contributes to the voltage fluctuations. The thermal JN noise (\( S_{Vb} = 4k_BT R \)) can be added to obtain the total voltage noise:
\[
S_V(\omega) = 4k_BT (R + R_1 \cos^2 \alpha). \tag{16}
\]
Our direct derivation of the PSD [Eqs. (15) and (16)] is consistent with the FDT combined with the angle-dependent resistance [15,16]. Thus, the present analysis can be considered an alternative derivation of the SMR effect.

IV. CONCLUSION

In summary, we report, to the best of our knowledge, the first observation of what we call spin Hall noise. The magnetization direction-dependent voltage noise and resistance measured in a YIG/Pt bilayer obey the FDT, confirming that spin Hall current-based physics of the SMR [15,16] implies the presence of the spin Hall noise. A theoretical description for the latter in terms of spin-current fluctuations gives insight into the nontrivial nature of entanglement of the spin contribution with the magnetization dynamics. In light of the FDT, observation of spin-current fluctuations emphasizes the dissipative nature of pure spin currents. The experimental resolution of the spin-current noise demonstrated here paves the way for advanced noise spectroscopy studies, such as (nonequilibrium) resistance [34] and spin-pumping shot noise.

ACKNOWLEDGMENTS


APPENDIX: SPIN DIFFUSION IN 3D

Here, we solve the spin-diffusion equation in the normal metal (N) and calculate the spin-current correlators required for evaluating the voltage noise power spectral density (PSD). We show that a three-dimensional analysis yields the same result as the quasi-one-dimensional model [Eq. (9)]. The notations and the coordinate system are defined in Fig. 3.

Since we are interested in the \( x \)-polarized component of the spin current \( \mathbf{J}_s^x(r,t) = -D \nabla \mu_s^x \) at frequencies much smaller than the inverse spin-flip rate, we have to solve the time-independent spin-diffusion equation:
\[
\nabla^2 \mu_s^x = \frac{\mu_s^x}{\lambda_{sd}}, \tag{A1}
\]
with the boundary conditions \( \mathbf{J}_s^x(\mathbf{r},t) = -D \nabla \mu_s^x = \mathbf{J}_s^x(\mathbf{q},t) \) [see Eq. (4)] at \( y = 0 \) and \( \mathbf{J}_s^x(\mathbf{r},t) = 0 \) at \( y = t_N \), where \( \mathbf{J}_s^x \) denotes the \( x \)-polarized spin current flowing along the \( y \) direction. This equation is valid for frequencies much smaller than the spin-flip rate (~THz in Pt). Physically, all time dependence comes from the boundary conditions to which the spin accumulation reacts instantaneously. The general solution for a translationally invariant planar system reads [32,35]
\[
\mu_s^x = \sum_k \frac{\tilde{\mu}_s^x(k)}{A} e^{i\mathbf{k} \cdot \mathbf{r}} \frac{\cos [\mathbf{k} \cdot \mathbf{t}_N]}{\sinh (\mathbf{k} \cdot \mathbf{t}_N)}, \tag{A2}
\]
where \( A \) is the interface area, \( \mathbf{k} \) an in-plane wave vector, \( c_k^x = \sqrt{1/\lambda_{sd} + k_x^2 + k_y^2} \), and spin current,
\[
\mathbf{J}_s^x(\mathbf{r},t) = \sum_k \frac{\tilde{J}_s^x(\mathbf{k},t)}{A} e^{i\mathbf{k} \cdot \mathbf{r}} f(\mathbf{k},y), \tag{A3}
\]
with
\[
f(\mathbf{k},y) = \frac{\sinh [c_k^x(t_N - y)]}{\sinh (c_k^x t_N)}. \tag{A4}
\]

The voltage autocorrelation and PSD are governed by the integral over the metal film:
\[
g(t) = \int \int \langle J_s^x(\mathbf{r},t) J_s^x(\mathbf{r}',0) \rangle d^3r d^3r', \tag{A5}
\]
\[
= \int \int F(\mathbf{r},\mathbf{r}',t) d^3r d^3r', \tag{A6}
\]
where \( \langle \cdot \rangle \) denotes statistical averaging. With Eq. (A3),
\[
F(\mathbf{r},\mathbf{r}',t) = \frac{1}{A^2} \sum_{k,k'} \langle \tilde{J}_s^x(\mathbf{k},t) \tilde{J}_s^x(\mathbf{k}',0) \rangle \times e^{i(\mathbf{k} \cdot \mathbf{r} + \mathbf{k}' \cdot \mathbf{r}')} f(\mathbf{k},y) f(\mathbf{k}',y'). \tag{A7}
\]
Due to the boundary condition at \( y = 0 \), \( \tilde{J}_s^x(\mathbf{k},t) \) is the Fourier transform of \( J_s^x(\mathbf{q},t) \), whence, employing Eq. (8),
\[
\langle \tilde{J}_s^x(\mathbf{k},t) \tilde{J}_s^x(\mathbf{k}',0) \rangle = \int \int \langle \tilde{J}_s^x(\mathbf{q},t) \tilde{J}_s^x(\mathbf{q}',0) \rangle e^{-i(\mathbf{k} \cdot \mathbf{q} + \mathbf{k}' \cdot \mathbf{q}')} d^2\mathbf{q} d^2\mathbf{q}', \tag{A8}
\]
\[
= \frac{\hbar g_{s}}{2\pi} k_BT A \delta_{\mathbf{k} - \mathbf{k}'} \delta(t), \tag{A9}
\]
and
\[
F(\mathbf{r},\mathbf{r}',t) = \frac{1}{A} \sum_k \frac{\hbar g_{s}}{2\pi} k_BT \delta(t) e^{i(\mathbf{k} \cdot \mathbf{r} - \mathbf{k}' \cdot \mathbf{r}')} f(\mathbf{k},y) f(\mathbf{k}',y'). \tag{A10}
\]
Therefore,

\[
g(t) = \int \int \frac{1}{A} \sum_k \frac{\hbar \tilde{g}_r}{2\pi} k_B T \delta(t) e^{ik \cdot \mathbf{q} - \phi} f(k, y) f(k', y') d^3r d^3r',
\]

(A11)

\[
eq \frac{\hbar \tilde{g}_r}{2\pi} k_B T A \delta(t) \left[ \int_0^{t_N} \frac{\sinh \left( \frac{t y}{l_N} \right)}{\sinh \left( \frac{t y}{l_N} \right)} dy \right]^2,
\]

(A12)

\[
eq \frac{\hbar \tilde{g}_r}{2\pi} k_B T A \lambda_{sd}^2 \tanh^2 \left( \frac{t N}{2\lambda_{sd}} \right) \delta(t),
\]

(A13)

which agrees with Eq. (12). The volume integral of the electromotive force that amounts to the total voltage across \( N \) corresponds to the \( k = 0 \) component of the in-plane variations, thereby reducing the 3D to an effectively 1D problem.