Letter

Motion-induced spin transfer

Daigo Oue [1,2,*] and Mamoru Matsuo [1,3,4,5]

[1] *Kavli Institute for Theoretical Sciences, University of Chinese Academy of Sciences, Beijing 100190, China*
[2] *The Blackett Laboratory, Department of Physics, Imperial College London, Prince Consort Road, Kensington, London SW7 2AZ, United Kingdom*
[3] *CAS Center for Excellence in Topological Quantum Computation, University of Chinese Academy of Sciences, Beijing 100190, China*
[4] *Advanced Science Research Center, Japan Atomic Energy Agency, Tokai 319-1195, Japan*
[5] *RIKEN Center for Emergent Matter Science (CEMS), Wako, Saitama 351-0198, Japan*

(Received 19 October 2021; revised 27 December 2021; accepted 3 January 2022; published 11 January 2022)

We propose a spin transport induced by inertial motion. Our system is composed of two host media and a narrow vacuum gap in between. One of the hosts is sliding at a constant speed relative to the other. This mechanical motion causes the Doppler effect, which shifts the density of states and the nonequilibrium distribution function in the moving medium. Those shifts induce the difference in the distribution function between the two media, and they result in tunneling spin current. The spin current is calculated from the Schwinger-Keldysh formalism with a spin tunneling Hamiltonian. This scheme does not require temperature difference, voltage, or chemical potential.

DOI: 10.1103/PhysRevB.105.L020302

Introduction. Transport is a universal phenomenon in physics. Utilizing electron, neutron, and photon transports in free space has provided high precision measurements such as microscopy and spectroscopy, which have played important roles not only in physics but also in materials science, chemistry, and biology. Precisely guiding those excitations in media has enabled electrical and optical communications and information storage. Recent advances in condensed-matter physics have realized not only the transport of a single quantum but also the manipulation of its properties.

In an emerging field called spintronics, manipulation of the spin angular momenta of electrons has been conducted in various ways. For example, spin tunneling transport at the interface between a normal metal and a ferromagnetic insulator can be driven by microwave irradiation on the ferromagnetic side. This type of spin transport is known as the spin pumping effect [1–5]. We proposed that visible light could also be used to drive spin transport at metallic interfaces [6–8]. There is another popular way to drive the spin tunneling making use of the temperature difference between two media, i.e., the spin Seebeck effect [9–12]. In these schemes, the differences in the nonequilibrium distributions between the two media drove the spin transports.

Another interesting direction in spintronics is to use a mechanical degree of freedom for spin manipulation. Since spin is a kind of angular momenta, it can be manipulated by mechanical rotation in accordance with the angular momentum conservation. Indeed, Barnett, Einstein, and de Haas showed experimentally that rigid-body rotation interacts with a magnetic moment originating from the spin angular momenta of electrons [13,14]. The mechanical manipulation of spin is demonstrated in a variety of systems, including micromechanical systems [15–18], microfluid systems [19–21], atomic nuclei [22,23], and quark-gluon plasma [24]. These effects can be comprehensively understood as consequences of spin-rotation coupling [25,26].

Although there are various studies on spin transport and manipulation by mechanical motion as reviewed in Ref. [27], there is no study on spin tunneling transport driven by mechanical motion.

In this work, we show that spin tunneling transport between two media can be induced by inertial motion. Our proposal is closely related to noncontact friction [28–31]. There are various theoretical works describing translational-type [32–38] and rotational-type [39–41] noncontact friction. From the source point of view, Langevin-type equations have been used to describe the frictional force in fluctuating fields [28–30,32,36]. On the other hand, from the field point of view, the spectral shift induced by motion plays a vital role [33–35,37]. The spectral shifts produce the photon momentum flux between two objects and hence the friction. In other words, the mechanical motion empowers the linear momenta to be transferred from one to the other. Here, we consider spin transfer between relatively moving media instead of linear momentum transfer.

When relative motion is forced in a system, the system becomes inhomogeneous, i.e., driven into a nonequilibrium state. We can tell one medium from the other, and there is no longer symmetry in the direction normal to the surfaces. This symmetry breaking induced by the nonuniformly forced motion provides a possibility of current generation in the direction. To take the inhomogeneity into consideration, we utilize the nonequilibrium (Schwinger-Keldysh) Green's function, and we perturbatively evaluate the effects of relative motion, e.g., spin currents.

*daigo.oue@gmail.com

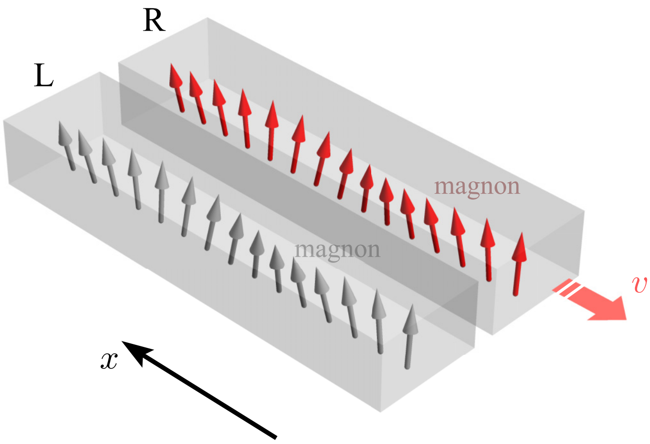

FIG. 1. The schematic image of the setup considered in this study. A very narrow gap separates two media hosting magnons. The right medium is moving in the $-x$ direction at a constant velocity v while the left medium stands still. Due to the Doppler effect, the dispersion relation of the magnon in the right medium observed in the laboratory frame is shifted.

We consider two media separated by a very narrow gap (Fig. 1).

Each medium hosts a magnon and is described by the following Hamiltonian within the Holstein-Primakoff approximation [42]:

$$H_0 = E_0 + \sum_k \hbar\omega_k b_k^\dagger b_k, \tag{1}$$

where E_0 is the classical ground-state energy of the medium, ω_k is the magnon dispersion relation, and we have introduced bosonic creation and annihilation operators.

The two media are interacting via the following tunneling Hamiltonian:

$$H_{\rm int} = \sum_k H_{\rm ex} b_{Rk} b_{Lk}^\dagger + {\rm H.c.}, \tag{2}$$

where $H_{\rm ex}$ is a coupling strength, and the subscripts L and R specify the medium. Although we assume for the sake of simplicity that the coupling strength $H_{\rm ex}$ is constant in this study, it could be dependent on the wave number k.

Spin current induced by inertial motion. Let us consider the change of the spin on the left medium in the interaction picture. We can obtain

$$\frac{\partial}{\partial t}\sum_k \langle S_k^z(t)\rangle = -\sum_k 2H_{\rm ex}\,{\rm Im}\langle b_{Rk}(t) b_{Lk}^\dagger(t)\rangle, \tag{3}$$

where $S_k^z = S - b_{Lk}^\dagger b_{Lk}$ is the z component of the spin in the left medium, and $\langle\cdots\rangle$ denotes the average with respect to the full Hamiltonian. Let us define the spin current flowing into the left medium at $t = t_1$,

$$\langle I_s(t_1)\rangle \equiv -\sum_k 2H_{\rm ex}\,{\rm Im}\langle b_{Rk}(t_1 - 0) b_{Lk}^\dagger(t_1)\rangle. \tag{4}$$

We shall omit $\sum_k$ in the following where relevant.

Using the formal perturbative expansion [43], we can evaluate the spin current up to the second order in the coupling strength $H_{\rm ex}$,

$$\langle I_s(t_1)\rangle = \frac{2H_{\rm ex}^2}{\hbar}{\rm Re}\int_C \langle \mathcal{T}_C b_{Lk}^\dagger(t_1^-) b_{Lk}(t_2)\rangle_0 \times \langle \mathcal{T}_C b_{Rk}(t_1^+) b_{Rk}^\dagger(t_2)\rangle_0 dt_2, \tag{5}$$

where $\langle\cdots\rangle_0$ is the average with respect to the unperturbed Hamiltonian, and $\mathcal{T}_C$ is the time-ordering operator on the Schwinger-Keldysh contour composed of a forward branch C_+ and a backward one C_- (see Fig. 2). Note that $t^\pm$ denotes time on the forward and backward branches.

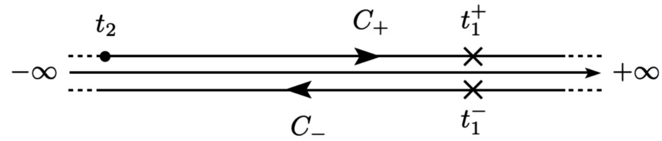

FIG. 2. Schwinger-Keldysh contour composed of forward and backward branches $C_\pm$.

Here, we introduce the nonequilibrium Green's function,

$$\chi_k(t_1, t_2) := \frac{1}{i\hbar}\langle \mathcal{T}_C b_k(t_1) b_k^\dagger(t_2)\rangle_0, \tag{6}$$

whose lesser and greater components read

$$\chi_{k;12}^< := \chi_k(t_1^+, t_2^-) = \frac{-i}{\hbar}\langle b_k^\dagger(t_2) b_k(t_1)\rangle_0, \tag{7}$$

$$\chi_{k;12}^> := \chi_k(t_1^-, t_2^+) = \frac{-i}{\hbar}\langle b_k(t_1) b_k^\dagger(t_2)\rangle_0. \tag{8}$$

We can write the chronologically ordered and antichronologically ordered components in terms of the lesser and greater components [(7),(8)],

$$\chi_{k;12}^{++} := \theta(t_1 - t_2)\chi_{k;12}^> + \theta(t_2 - t_1)\chi_{k;12}^<, \tag{9}$$

$$\chi_{k;12}^{--} := \theta(t_1 - t_2)\chi_{k;12}^< + \theta(t_2 - t_1)\chi_{k;12}^>, \tag{10}$$

where θ denotes the Heaviside unit step function.

Splitting the contour C into the forward and backwards parts, we can write the real-time representation,

$$\frac{\langle I_s(t_1)\rangle}{2\hbar H_{\rm ex}^2} = -{\rm Re}\int \left(\chi_{Rk;12}^{\mathfrak{R}}\chi_{Lk;21}^< + \chi_{Rk;12}^<\chi_{Lk;21}^{\mathfrak{A}}\right) dt_2, \tag{11}$$

where we have defined the retarded and advanced components,

$$\chi_{k;12}^{\mathfrak{R}} := \frac{-i}{\hbar}\theta(t_1 - t_2)\langle [b_k(t_1), b_k^\dagger(t_2)]\rangle_0 = \chi_{k;12}^{++} - \chi_{k;12}^<, \tag{12}$$

$$\chi_{k;12}^{\mathfrak{A}} := \frac{+i}{\hbar}\theta(t_2 - t_1)\langle [b_k(t_1), b_k^\dagger(t_2)]\rangle_0 = \chi_{k;12}^< - \chi_{k;12}^{--}, \tag{13}$$

which is merely the dynamical (magnetic) susceptibility of the medium. Note that the square brackets denotes the commutation relation here, i.e., $[\bullet, \circ] = \bullet\circ - \circ\bullet$.

In the steady state, Green's functions depend only on the time difference, e.g.,

$$\chi_{k;12}^{\mathfrak{R}} = \frac{1}{2\pi}\int \chi_{k\omega}^{\mathfrak{R}} e^{-i\omega(t_1 - t_2)} d\omega. \tag{14}$$

Thus, working on the frequency domain, we can simplify the integral (11) in the steady state,

$$\langle I_s^{ss} \rangle = \frac{4H_{ex}^2}{2\pi} \sum_{k>0} \int \Delta j_s^{ss}(k,\omega) d\omega, \tag{15}$$

$$j_s^{ss}(k,\omega) = \hbar \,\mathrm{Im}\, \chi_{Lk\omega}^{\mathfrak{R}} \,\mathrm{Im}\, \chi_{Rk\omega}^{\mathfrak{R}} \delta n_k. \tag{16}$$

Here, we have symmetrized the integrand with respect to the wave number, $\Delta j_s^{ss}(k,\omega) = j_s^{ss}(k,\omega) + j_s^{ss}(-k,\omega)$, and we defined the distribution difference between the two media, $\delta n_k := n_b(\omega_{Lk}) - n_b(\omega_{Rk})$, where n_b denotes the Bose distribution function, and $\omega_{L(R)k}$ is the magnon dispersion in the left (right) medium. Note that we have used the Kadanoff-Baym ansatz, i.e., $\chi_{k\omega}^{<} = 2in_b(\omega_k)\,\mathrm{Im}\,\chi_{k\omega}^{\mathfrak{R}}$, in order to get Eq. (15).

That is the formula we use to evaluate the spin current flowing between the two media. The integrand $\Delta j_s^{ss}(k,\omega)$ is composed of (i) the products of magnon spectra $\mathrm{Im}\,\chi_{Lk\omega}^{\mathfrak{R}}\,\mathrm{Im}\,\chi_{Rk\omega}^{\mathfrak{R}}$ and (ii) the distribution difference δn_k. This implies that large spectral overlap and large population differences between the two media drive large spin currents.

In our setup, the inertial motion of the right medium is the key to generate a finite population difference. To consider the effects of inertial motion, we shall define the physical quantities in the comoving frame and go back to the laboratory frame. In the nonrelativistic regime ($|v/c| \ll 1$), the Lorentz boost can be safely approximated by the Galilean boost, which we use to go back to the laboratory frame,

$$t \to t, \quad x \to x - vt. \tag{17}$$

Applying the Galilean boost to a function $\psi(x)$, we have

$$\psi(x) = \frac{1}{(2\pi)^2} \iint \psi_{k\omega} e^{i(kx-\omega t)} dk\, d\omega, \tag{18}$$

$$\to \frac{1}{(2\pi)^2} \iint \psi_{k,\omega+vk} e^{i(kx-\omega t)} dk\, d\omega. \tag{19}$$

This implies that the Galilean boost (17) induces the Doppler effect, i.e., the spectra and hence the dispersion relations in moving media are shifted ($\omega_k \to \omega_k - vk$).

In our case, the spectrum and the distribution function of magnons in the right medium are shifted,

$$\mathrm{Im}\,\chi_{Rk\omega}^{\mathfrak{R}} \to \mathrm{Im}\,\chi_{Rk,\omega+vk}^{\mathfrak{R}}, \quad n_b(\omega_{Rk}) \to n_b(\omega_{Rk} - vk). \tag{20}$$

When the two media are made of the same material, we substitute

$$\mathrm{Im}\,\chi_{Lk\omega}^{\mathfrak{R}} = \mathrm{Im}\,\chi_{k\omega}^{\mathfrak{R}}, \quad \mathrm{Im}\,\chi_{Rk\omega}^{\mathfrak{R}} = \mathrm{Im}\,\chi_{k,\omega+vk}^{\mathfrak{R}}, \quad \delta n_k = n_b(\omega_k) - n_b(\omega_k - vk). \tag{21}$$

From the expression δn_k, we can immediately find that there is no spin current if the right medium is not moving ($v = 0$). In the following, we assume a simple parabolic dispersion for the magnon, $\omega_k = Dk^2 + \omega_0$, where $\omega_0 = \gamma B$ is the Zeeman energy. The retarded component of the magnon Green's function can be given in the frequency domain,

$$\chi_{k\omega}^{\mathfrak{R}} = \frac{1/\hbar}{\omega - \omega_k + i\Gamma}, \tag{22}$$

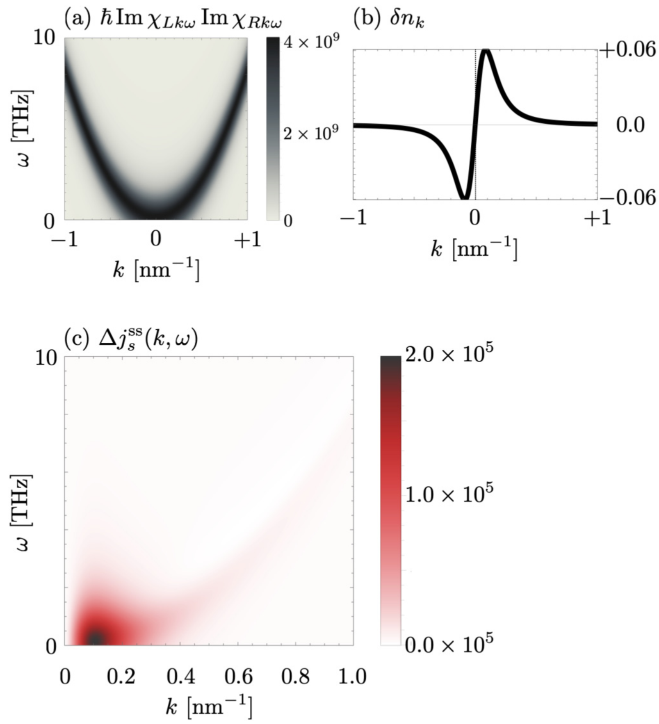

FIG. 3. (a) The spectral overlap of magnons in left and right media $\mathrm{Im}\,\chi_{Lk\omega}^{\mathfrak{R}}\,\mathrm{Im}\,\chi_{Rk\omega}^{\mathfrak{R}}$. (b) Magnon distribution difference between left and right media δn_k. (c) The integrand $\Delta j_s^{ss}(k,\omega)$ that yields the tunneling spin current. We set the spectral broadening due to impurities, etc. $\Gamma = 1$ (meV) ≈ 0.24 (THz), the velocity of the right medium $v = 1$ (m s^{-1}), the static magnetic field $B = 1$ (T), and $D = 532$ (meV Å^2) in accordance with Ref. [44].

where Γ is spectral broadening, for example, due to surface roughness and impurity scattering.

Note that we can analyze the spin current in a frame comoving with the right medium, and the result does not contradict the calculation in the laboratory frame [43]. In Fig. 3(c), the integrand $\Delta j_s^{ss}(k,\omega)$ in the spin current formula (15) is plotted. Since the Doppler shift $\Delta\omega_k = vk$ is smallest when $k = 0$, the spectral overlap $\mathrm{Im}\,\chi_{Lk\omega}^{\mathfrak{R}}\,\mathrm{Im}\,\chi_{Rk\omega}^{\mathfrak{R}}$ becomes large in the low-frequency region [see Fig. 3(a)]. We can see that the amount of the distribution difference δn_k is large in the low-wave-number region [see Fig. 3(b)]. This implies that the dominant contribution to the steady-state spin current $\langle I_s^{ss}\rangle$ comes from that region and justifies introducing the cutoff frequency and the wave number when evaluating the integral (15) in numerics. Numerically integrating the spin current formula (15), we can obtain Figs. 4 and 5. To obtain those figures, we substituted a far smaller number into the coupling strength than the magnon energy ($H_{ex} \ll \hbar\omega_0 < \hbar\omega_k$), and thus we can safely adopt the perturbative evaluation of the spin current.

We have plotted the spin current $\langle I_s^{ss}\rangle$ as a function of the right medium velocity v for three different temperatures in Fig. 4. As the velocity v is larger, the Doppler shift $\Delta\omega = vk$, and hence the distribution difference δn_k increases. This is why the spin current increases with the velocity of the right medium. We can fit the spin current by a parabolic function.

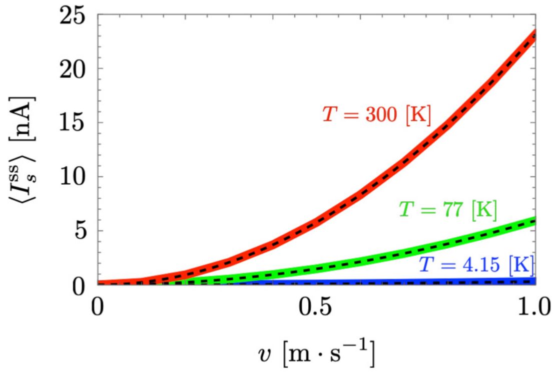

FIG. 4. Spin current at the steady state $\langle I_s^{ss} \rangle$ as a function of the velocity of the right medium v. Blue, gray, and red curves are the spin currents at liquid helium, liquid nitrogen, and room temperatures, [$T = 4.15, 77, 300$ (K)]. We set the coupling strength $H_{ex} = 1$ (GHz) ≈ 50 (meV), which is far smaller compared with the magnon frequency (i.e., $H_{ex} \ll \hbar\omega_k$). The other parameters are the same as the previous figure. Each curve can be fitted by a parabolic function (black dashed curve). This implies that the motion-induced spin transfer is the second-order effect.

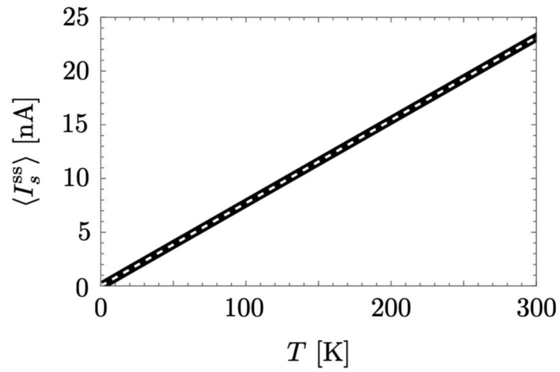

FIG. 5. Spin current at the steady state $\langle I_s^{ss} \rangle$ as a function of temperature T. Since the spin current can be fitted by a linear function (white dashed line), the motion-induced spin transfer is proportional to the temperature T. We set $v = 1$ (m s^{-1}) and other parameters the same as the previous figures.

This reflects the fact that the leading term of the motion-induced spin current is the second order, i.e., $\langle I_s^{ss} \rangle \propto H_{ex}^2$.

In Fig. 5, we show the temperature dependence of the spin current. The spin current can be linearly fitted (see the white line in the figure) and is proportional to the temperature T.

Conclusions. In this Letter, we proposed motion-induced spin transfer between two media which host and can exchange magnons. One of the two media is moving at a constant velocity, and the inertial motion causes the Doppler effect. This results in the spectral shift of the magnon spectrum and distribution function in the moving medium. According to our perturbative calculation within the Schwinger-Keldysh formalism, the difference in the magnon distribution between the two media drives the spin transfer from the moving one to the other.

As for the possibility of the experimental verification of our proposal, we could use the state-of-the-art inverse spin Hall measurement that was used, for example, in Ref. [19], which can detect an electric signal of the order of 1 nV.

Our proposal will open a new door to spin manipulation by inertial motion.

Acknowledgments. D.O. and M.M. deeply thank Yuya Ominato for fruitful discussion on the nonequilibrium Green's function method. D.O. is funded by the President's Ph.D. Scholarships at Imperial College London. This work is partially supported by the Priority Program of Chinese Academy of Sciences, Grant No. XDB28000000 and a Grant-in-Aid for Scientific Research B (20H01863) from MEXT, Japan.

[1] R. H. Silsbee, A. Janossy, and P. Monod, Coupling between ferromagnetic and conduction-spin-resonance modes at a ferromagnetic—normal-metal interface, Phys. Rev. B **19**, 4382 (1979).
[2] Y. Tserkovnyak, A. Brataas, and G. E. W. Bauer, Enhanced Gilbert Damping in Thin Ferromagnetic Films, Phys. Rev. Lett. **88**, 117601 (2002).
[3] I. Žutić, J. Fabian, and S. D. Sarma, Spintronics: Fundamentals and applications, Rev. Mod. Phys. **76**, 323 (2004).
[4] Y. Ohnuma, H. Adachi, E. Saitoh, and S. Maekawa, Enhanced dc spin pumping into a fluctuating ferromagnet near T_C, Phys. Rev. B **89**, 174417 (2014).
[5] T. Kato, Y. Ohnuma, M. Matsuo, J. Rech, T. Jonckheere, and T. Martin, Microscopic theory of spin transport at the interface between a superconductor and a ferromagnetic insulator, Phys. Rev. B **99**, 144411 (2019).
[6] D. Oue and M. Matsuo, Electron spin transport driven by surface plasmon polaritons, Phys. Rev. B **101**, 161404(R) (2020).
[7] D. Oue and M. Matsuo, Effects of surface plasmons on spin currents in a thin film system, New J. Phys. **22**, 033040 (2020).
[8] D. Oue and M. Matsuo, Optically induced electron spin currents in the Kretschmann configuration, Phys. Rev. B **102**, 125431 (2020).
[9] K. Uchida, S. Takahashi, K. Harii, J. Ieda, W. Koshibae, K. Ando, S. Maekawa, and E. Saitoh, Observation of the spin Seebeck effect, Nature (London) **455**, 778 (2008).
[10] K.-i. Uchida, J. Xiao, H. Adachi, J.-i. Ohe, S. Takahashi, J. Ieda, T. Ota, Y. Kajiwara, H. Umezawa, H. Kawai *et al.*, Spin Seebeck insulator, Nat. Mater. **9**, 894 (2010).
[11] K.-i. Uchida, H. Adachi, T. Ota, H. Nakayama, S. Maekawa, and E. Saitoh, Observation of longitudinal spin-Seebeck effect in magnetic insulators, Appl. Phys. Lett. **97**, 172505 (2010).
[12] H. Adachi, J.-i. Ohe, S. Takahashi, and S. Maekawa, Linear-response theory of spin Seebeck effect in ferromagnetic insulators, Phys. Rev. B **83**, 094410 (2011).
[13] A. Einstein and W. J. de Haas, Experimental proof of the existence of Ampère's molecular currents, Deutsche Physikalische Gesellschaft, Verhandlungen **17**, 152 (1915).
[14] S. J. Barnett, Magnetization by rotation, Phys. Rev. **6**, 239 (1915).

[15] T. M. Wallis, J. Moreland, and P. Kabos, Einstein–de Haas effect in a NiFe film deposited on a microcantilever, Appl. Phys. Lett. **89**, 122502 (2006).
[16] G. Zolfagharkhani, A. Gaidarzhy, P. Degiovanni, S. Kettemann, P. Fulde, and P. Mohanty, Nanomechanical detection of itinerant electron spin flip, Nat. Nanotechnol. **3**, 720 (2008).
[17] D. Kobayashi, T. Yoshikawa, M. Matsuo, R. Iguchi, S. Maekawa, E. Saitoh, and Y. Nozaki, Spin Current Generation Using a Surface Acoustic Wave Generated via Spin-Rotation Coupling, Phys. Rev. Lett. **119**, 077202 (2017).
[18] K. Harii, Y.-J. Seo, Y. Tsutsumi, H. Chudo, K. Oyanagi, M. Matsuo, Y. Shiomi, T. Ono, S. Maekawa, and E. Saitoh, Spin Seebeck mechanical force, Nat. Commun. **10**, 2616 (2019).
[19] R. Takahashi, M. Matsuo, M. Ono, K. Harii, H. Chudo, S. Okayasu, J. Ieda, S. Takahashi, S. Maekawa, and E. Saitoh, Spin hydrodynamic generation, Nat. Phys. **12**, 52 (2016).
[20] H. Tabaei Kazerooni, A. Thieme, J. Schumacher, and C. Cierpka, Electron spin-vorticity coupling in pipe flows at low and high Reynolds number, Phys. Rev. Appl. **14**, 014002 (2020).
[21] H. Tabaei Kazerooni, G. Zinchenko, J. Schumacher, and C. Cierpka, Electrical voltage by electron spin-vorticity coupling in laminar ducts, Phys. Rev. Fluids **6**, 043703 (2021).
[22] H. Chudo, M. Ono, K. Harii, M. Matsuo, J. Ieda, R. Haruki, S. Okayasu, S. Maekawa, H. Yasuoka, and E. Saitoh, Observation of Barnett fields in solids by nuclear magnetic resonance, Appl. Phys. Express **7**, 063004 (2014).
[23] A. Wood, E. Lilette, Y. Fein, V. Perunicic, L. Hollenberg, R. Scholten, and A. Martin, Magnetic pseudo-fields in a rotating electron–nuclear spin system, Nat. Phys. **13**, 1070 (2017).
[24] L. Adamczyk, J. Adkins, G. Agakishiev, M. Aggarwal, Z. Ahammed, N. Ajitanand, I. Alekseev, D. Anderson, R. Aoyama, A. Aparin *et al.*, Global λ hyperon polarization in nuclear collisions, Nature (London) **548**, 62 (2017).
[25] M. Matsuo, J. Ieda, K. Harii, E. Saitoh, and S. Maekawa, Mechanical generation of spin current by spin-rotation coupling, Phys. Rev. B **87**, 180402(R) (2013).
[26] M. Matsuo, Y. Ohnuma, and S. Maekawa, Theory of spin hydrodynamic generation, Phys. Rev. B **96**, 020401(R) (2017).
[27] M. Matsuo, E. Saitoh, and S. Maekawa, Spin-mechatronics, J. Phys. Soc. Jpn. **86**, 011011 (2017).
[28] I. Dorofeyev, H. Fuchs, G. Wenning, and B. Gotsmann, Brownian Motion of Microscopic Solids under the Action of Fluctuating Electromagnetic Fields, Phys. Rev. Lett. **83**, 2402 (1999).
[29] B. C. Stipe, H. J. Mamin, T. D. Stowe, T. W. Kenny, and D. Rugar, Noncontact Friction and Force Fluctuations between Closely Spaced Bodies, Phys. Rev. Lett. **87**, 096801 (2001).
[30] B. Gotsmann and H. Fuchs, Dynamic Force Spectroscopy of Conservative and Dissipative Forces in an Al-Au (111) Tip-Sample System, Phys. Rev. Lett. **86**, 2597 (2001).
[31] K. Saitoh, K. Hayashi, Y. Shibayama, and K. Shirahama, Gigantic Maximum of Nanoscale Noncontact Friction, Phys. Rev. Lett. **105**, 236103 (2010).
[32] J. R. Zurita-Sánchez, J.-J. Greffet, and L. Novotny, Friction forces arising from fluctuating thermal fields, Phys. Rev. A **69**, 022902 (2004).
[33] A. I. Volokitin and B. N. J. Persson, Near-field radiative heat transfer and noncontact friction, Rev. Mod. Phys. **79**, 1291 (2007).
[34] A. I. Volokitin and B. N. J. Persson, Quantum Friction, Phys. Rev. Lett. **106**, 094502 (2011).
[35] M. G. Silveirinha, Theory of quantum friction, New J. Phys. **16**, 063011 (2014).
[36] J.-H. She and A. V. Balatsky, Noncontact Friction and Relaxational Dynamics of Surface Defects, Phys. Rev. Lett. **108**, 136101 (2012).
[37] K. Milton, J. Høye, and I. Brevik, The reality of Casimir friction, Symmetry **8**, 29 (2016).
[38] S. M. Barnett and M. Sonnleitner, Vacuum friction, J. Mod. Opt. **65**, 706 (2018).
[39] A. Manjavacas and F. J. Garcia de Abajo, Vacuum Friction in Rotating Particles, Phys. Rev. Lett. **105**, 113601 (2010).
[40] R. Zhao, A. Manjavacas, F. J. Garcia de Abajo, and J. B. Pendry, Rotational Quantum Friction, Phys. Rev. Lett. **109**, 123604 (2012).
[41] F. Intravaia, M. Oelschläger, D. Reiche, D. A. R. Dalvit, and K. Busch, Quantum Rolling Friction, Phys. Rev. Lett. **123**, 120401 (2019).
[42] T. Holstein and H. Primakoff, Field dependence of the intrinsic domain magnetization of a ferromagnet, Phys. Rev. **58**, 1098 (1940).
[43] See Supplemental Material at http://link.aps.org/supplemental/10.1103/PhysRevB.105.L020302 for an evaluation of the spin current using the formal perturbative expansion.
[44] A. J. Princep, R. A. Ewings, S. Ward, S. Tóth, C. Dubs, D. Prabhakaran, and A. T. Boothroyd, The full magnon spectrum of yttrium iron garnet, npj Quantum Mater. **2**, 63 (2017).