Selective Activation of an Isolated Magnetic Skyrmion in a Ferromagnet with Microwave Electric Fields

Akihito Takeuchi\textsuperscript{1,\textcircled{a}} and Masahito Mochizuki\textsuperscript{2,3,\textcircled{b}}
\textsuperscript{1) Department of Physics and Mathematics, Aoyama Gakuin University, Sagamihara, Kanagawa 252-5258, Japan
\textsuperscript{2) Department of Applied Physics, Waseda University, Okubo, Shinjuku-ku, Tokyo 169-8555, Japan
\textsuperscript{3) PRESTO, Japan Science and Technology Agency, Kawaguchi, Saitama 332-0012, Japan}

We theoretically reveal that pure eigenmodes of an isolated magnetic skyrmion embedded in a ferromagnetic environment can be selectively activated using microwave electric fields without exciting gigantic ferromagnetic resonances, in contrast to conventional methods using microwave magnetic fields. We also demonstrate that this selective activation of a skyrmion can efficiently drive its translational motion in a ferromagnetic nanotrack under application of an external magnetic field inclined from the normal direction. We find that a mode with combined breathing and rotational oscillations induces much faster skyrmion propagation than the breathing mode studied previously by Wang et al. [Phys. Rev. B \textbf{92}, 020403(R) (2015)].

PACS numbers: 76.50.+g,78.20.Ls,78.20.Bh,78.70.Gq

A skyrmion crystal, a hexagonally crystalized state of magnetic skyrmions\textsuperscript{11,16} has characteristic resonance modes at microwave frequencies\textsuperscript{7,13} which give rise to intriguing physical phenomena\textsuperscript{11,13} such as microwave directionnal dichroism\textsuperscript{12,13}, spin-voltage induction\textsuperscript{15,17}, and spin-current generation\textsuperscript{18}. When a static magnetic field $H_{\text{ex}}$ is applied perpendicular to a thin-plate specimen of the skyrmion-hosting magnet, several types of spin-wave modes emerge depending on the microwave polarization\textsuperscript{19}. A microwave magnetic field $H_{\omega}$ normal to the skyrmion plane ($H_{\parallel}$) activates the so-called breathing mode in which all the skyrmions constituting the skyrmion crystal uniformly expand and shrink in an oscillatory manner. On the other hand, the $H_{\omega}$ field within the skyrmion plane ($H_{\perp \omega}$) activates two types of rotation modes with opposite rotational senses, in which cores of all the skyrmions circulate uniformly in counterclockwise and clockwise ways.

In addition to the crystallized form, skyrmions can appear as individual defects in a ferromagnetic state; such skyrmions are also expected to have peculiar collective modes\textsuperscript{19}. Isolated skyrmions confined in a nanoferromagnet are potentially useful for applications\textsuperscript{20} to memory device\textsuperscript{21}, magnonics device\textsuperscript{22,23}, spin-torque oscillators\textsuperscript{24,25}, and microwave sensing device\textsuperscript{26}. As such, it is necessary to clarify the microwave-active eigenmodes of a single skyrmion in a ferromagnetic specimen.

In addition, it is important to establish a way to manipulate isolated skyrmions using microwaves for their device application. As the microwave field $H_{\perp \omega}$ cannot activate precession of the magnetizations when the microwave field is applied parallel to them, we can activate pure breathing-type skyrmion oscillations with $H_{\parallel \omega}$ under a perpendicular $H_{\text{ex}}$ field without exciting the background ferromagnetic state. However, once the $H_{\text{ex}}$ field is inclined, the microwave magnetic field $H_{\perp \omega}$ excites huge ferromagnetic resonances, which drown out the weaker skyrmion resonances. Moreover, the microwave energy is absorbed by the sample when exciting the gigantic ferromagnetic resonances, which would inevitably result in high energy consumption and considerable temperature rise. Therefore, a technique to selectively activate an isolated skyrmion is urgently required.

In this Letter, we first theoretically show that the eigenmodes of an isolated skyrmion embedded in a ferromagnetic environment can be selectively activated with a microwave electric field $E_{\omega}$ via oscillatory modulation of the Dzyaloshinskii-Moriya interaction (DMI) without ac-

\textbf{FIG. 1.} (color online). (a) Schematics of a magnetic bilayer system hosting skyrmions stabilized by the interfacial Dzyaloshinskii-Moriya interaction. (b) External magnetic field $H_{\text{ex}}$ where $\theta$ is an inclination angle from the normal direction. (c), (d) Color maps of the normal component of magnetizations $m_z$ (c) and scalar spin chiralities $c_s$ (d) of a skyrmion under a perpendicular $H_{\text{ex}}$ field. In-plane components of the magnetizations $(m_x,m_y)$ are displayed by arrows. (e), (f) Those under an inclined $H_{\text{ex}}$ field with $\theta=30^\circ$.

\textsuperscript{a)} Electronic mail: akihito@phys.aoyama.ac.jp
\textsuperscript{b)} Electronic mail: masa_mochizuki@waseda.jp
tivating ferromagnetic resonances. We then demonstrate that translational motion of the skyrmion can be driven through activating its resonance modes using this electric technique in an inclined $H_{\text{ex}}$ field. The latter part of the research was motivated by the recent theoretical work by Wang et al. that demonstrated skyrmion propagation via activating the breathing mode with a microwave $H_{\text{y}}$ field. Our study reveals that skyrmion motion can be driven not only by the breathing mode but also by other $E^{\omega}$-active modes. Furthermore, we find that a mode with combined clockwise and breathing oscillations can induce much faster propagation of the skyrmion than the previously studied breathing mode. Our findings pave a new way toward efficient manipulation of isolated skyrmions in nano-devices via the application of microwaves.

We consider a magnetic bilayer system composed of a ferromagnetic layer and a heavy-metal layer with strong spin-orbit interactions [see Fig. 1(a)], where the spatial inversion symmetry is broken at their interface, and thereby the DMI is active. An inclined magnetic field $H_{\text{ex}} = (H_x, 0, H_z)$ with $H_x = H_z \tan \theta$ is applied where $\theta$ is the inclination angle [see Fig. 1(b)]. For $0^\circ < \theta \leq 90^\circ$, the $H_{\text{ex}}$ field is inclined toward the positive $z$ direction. The DMI favors a rotating alignment of the magnetizations, which results in the formation of a Neel-type skyrmion. The skyrmion has a circular symmetry under a perpendicular $H_{\text{ex}}$ field ($\theta = 0^\circ$), but has disproportionate weight in distributions of the magnetizations and scalar spin chiralities [see Fig. 1(c) and (d)]. To describe the magnetism in this magnetic bilayer system, we employ a classical Heisenberg model on a square lattice with magnetization vectors $\mathbf{m}_i$ whose norm $m$ is unity. The Hamiltonian contains the ferromagnetic-exchange interaction, the Zeeman coupling to the magnetic fields, and the interfacial DMI:

$$H = -J \sum_{<i,j>} \mathbf{m}_i \cdot \mathbf{m}_j - [H_{\text{ex}} + H(t)] \cdot \sum_i \mathbf{m}_i + D(t) \sum_i [(\mathbf{m}_i \times \mathbf{m}_{i+\hat{x}}) \cdot \hat{y} - (\mathbf{m}_i \times \mathbf{m}_{i+\hat{y}}) \cdot \hat{x}].$$  

(1)

Here $H(t) = (0, 0, H_z(t))$ and $E(t) = (0, 0, E_z(t))$ represent time-dependent magnetic and electric fields applied perpendicular to the sample plane, respectively. We neglect magnetic anisotropies because they do not alter the results qualitatively although stability of the skyrmions and resonant frequencies of the eigenmodes may be slightly changed. The strength of the interfacial DMI can be tuned by applying a gate electric field normal to the plane via varying the extent of the spatial inversion asymmetry. The DMI coefficient $D(t) = D_0 + \Delta D(t)$ has two components, namely, a steady component $D_0$ and a $E(t)$-dependent component $\Delta D(t) = \kappa E_z(t)$ with $\kappa$ being the coupling constant. We take $J = 1$ for the energy units and take $D_0/J = 0.27$. For the inclined magnetic field, we take $H_z = 0.057$ with $\theta$ being a variable. The unit conversions when $J = 1$ meV are summarized in Table II.

| Exchange int. | $J = 1$ | 1 meV |
|---------------|---------|-------|
| Time          | $t = 1$ | 0.66 ps |
| Frequency $f = \omega/2\pi$ | $\omega = 0.01$ | 2.41 GHz |
| Magnetic field| $H = 1$ | 8.64 T |

TABLE I. Unit conversion table when $J = 1$ meV.

We simulate the magnetization dynamics by numerically solving the Landau-Lifshitz-Gilbert equation using the fourth-order Runge-Kutta method. The equation is given by,

$$\frac{d\mathbf{m}_i}{dt} = -\gamma m_i \times H_{\text{eff}}^{i} + \frac{\alpha_G}{m} m_i \times \frac{d\mathbf{m}_i}{dt}. \tag{2}$$

Here $\alpha_G = 0.04$ and $\gamma m$ are the Gilbert-damping constant and magnetic susceptibilities $\chi^{em}$ and $\chi^{mn}$.

$$\chi^{em}(\omega) = \sqrt{\frac{\mu_0}{\epsilon_0 \omega}} \Delta M^z_{\text{eff}}, \quad \chi^{mn}(\omega) = \frac{\Delta M^z_{\text{eff}}}{\mu_0 \omega}. \tag{3}$$

After applying a short pulse $H_z(t)$ or $E_z(t)$ with duration of $\Delta t = 1$ in the units of $1/\hbar$, we trace time profiles of the net magnetization $M_z(t) = (1/N) \sum_i m_{z,i}(t)$ and $M_{\text{osc}}(t) = M_z(t) - M_0(0)$ and obtain the Fourier transform $\Delta M^z_{\text{eff}}$. Dividing this quantity by an amplitude of the pulse $H_{\text{pulse}}$ or $E_{\text{pulse}}$, we obtain these susceptibilities. The calculations are performed using a system of $N = 160 \times 160$ sites with periodic boundary conditions.

Figure 2(a) displays the imaginary parts of the electromagnetic susceptibilities $\text{Im} \chi^{em}(\omega)$ for several values of $\theta$, which describe the response of the magnetizations to the $E^{\omega}$ field. When $\theta = 0^\circ$, the spectrum exhibits a single peak corresponding to the breathing mode activated by the oscillating DMI under the application of an AC electric field. When the $H_{\text{ex}}$ field is inclined with $\theta \neq 0^\circ$, two novel modes emerge, one with a higher and one with a lower frequency than the breathing mode. The intensities of the novel modes increase, whereas the intensity of the original breathing mode is increasingly suppressed as $\theta$ increases.

The imaginary parts of the magnetic susceptibilities $\text{Im} \chi^{mn}(\omega)$ in Fig. 2(b) describe the response of the magnetizations to the $H^{\omega}$ field. We find that only a breathing mode (m-Mode 2) appears when the $H_{\text{ex}}$ field is perpendicular ($\theta = 0^\circ$); however, its intensity decreases as $\theta$ increases. Remarkably, a large ferromagnetic resonance from the surrounding ferromagnetic magnetizations emerges under an inclined $H_{\text{ex}}$ field, whereas it is totally silent under the perpendicular $H_{\text{ex}}$ field. We also find a novel mode (m-Mode 3) at lower frequencies.

In reality, the skyrmion has another $H^{\omega}$-active mode (m-Mode 3) at higher frequencies, but it is hidden behind the gigantic ferromagnetic resonance and thus cannot be seen in the spectra of $\text{Im} \chi^{mn}(\omega)$. We can see this
FIG. 2. (color online). Imaginary parts of (a) the electromagnetic susceptibility Imχ^{em}(ω), (b) the magnetic susceptibility Imχ^{mm}(ω), and (c) the chirality susceptibility Imχ^{mc}(ω) for several values of $\theta$. Here, an inclined magnetic field $H_{ex} = (H_x \tan \theta, 0, H_z)$ with $H_z=0.057$ is applied. Three $E^\omega$-active modes are labeled as e-Modes 1-3 in (a), whereas the three $H^\omega$-active modes are labeled as m-Modes 1-3 in (b) and (c). The extremely intense mode around $\omega \sim 0.05 - 0.06$ in (b) is the ferromagnetic resonance (FMR).

A recent experiment for Ta/FeCoB/TaO$_x$ reported a huge $E$-field-induced variation of the interfacial DMI that reaches 140% when the applied voltage is 10 V. This observation supports the experimental feasibility of the 5% modula-
FIG. 4. (color online). (a) Illustration of skyrmion propagation driven by the $E^z$ field through activating a resonant mode of the skyrmion under an inclined $H_{ax}$ field. (b) Simulated snapshots of the skyrmion propagation for the $E^z$-active mode with $\omega=0.0513$ (m-Mode 2). (c)-(e) Trajectories (c) and drift velocities $v_x$ (d) and $v_y$ (e) of the propagating skyrmion driven by the $E^z$ field for three different resonant modes. (f)-(h) Those of the propagating skyrmion driven by the $H^z$ field. All the simulations were performed with an inclined $H_{ax}$ field where $H_z=0.057$ and $\theta=30^\circ$.

...tion assumed here. Figure 4(b) shows simulated snapshots of the skyrmion motion when the $E^z$ field with $\omega=0.0513$ is applied, which indeed displays propagation in the negative $x$ direction.

The trajectories of the propagating skyrmion during a time period from $t=0$ to $t=5000$ are shown in Fig. 4(c) for three different $E^z$-active modes at $\omega=0.0453, 0.0513,$ and $0.0664$. They were obtained by tracing the center-of-mass coordinate $(j_x, j_y)$ of the topological-charge distribution:

\[
j_{\gamma} = \sum_{i=(x,y)} c_{\gamma}(i_i, i_y) / \sum_{i=(x,y)} c_{\gamma}(i_i, i_y)
\]

with

\[
c_{\gamma} = \frac{1}{8\pi} \left[ m_i \cdot (m_{i+\hat{x}} \times m_{i+\hat{y}}) + m_i \cdot (m_{i-\hat{x}} \times m_{i-\hat{y}}) \right].
\]

We find that the direction and velocity of the motion vary depending on the skyrmion resonance mode. For all the modes, the skyrmion moves mainly in the negative $x$ direction. However, the trajectories for e-Mode 2 and 3 are slightly slanted toward the negative $y$ direction; meanwhile, the trajectory for e-Mode 1 is perfectly parallel to the $x$ axis. Interestingly, despite the slanted trajectory for e-Mode 2, its traveling distance along the $x$ axis is identical to that of the trajectory for e-Mode 1, which is directed perfectly along the $x$ axis. It can also be seen that the directions of the skyrmion movement for e-Mode 2 and 3 are the same, even though their traveling distances are different. The traveling distance for e-Mode 3 is much shorter than that for e-Mode 2 because of the smaller intensity of the latter mode, as can be seen in the inset of Fig. 4(c).

In Fig. 4(d) and (e), we plot the velocities $v=(v_x, v_y)$ of the skyrmion for three different $E^z$-active modes (see the right vertical axes) as functions of the strength of the time-dependent DMI, i.e., $\kappa E^z_x$. The values are calculated from the simulated displacements of the skyrmion in the $x$ and $y$ directions for a time period from $t=2000$ to $t=10000$ (see the left vertical axes) by assuming $J=1$ meV and $a=5\ \text{Å}$ with $a$ being the lattice constant. It can be seen that the velocities are proportional to the square of $\kappa E^z_x$, and they are on the order of $10^{-1} \text{ m/s}$. In fact, the traveling speed of the skyrmions achieved using this technique turns out to be relatively slow compared with the speed achieved in techniques based on electric-current injection. However, the present technique has an advantage: it is free from the Joule heating, and thus the energy consumption and the temperature rise could be significantly suppressed.

We finally study the skyrmion motion driven by AC magnetic fields $H^z$ under an inclined $H_{ax}$ field with $H_z=0.057$ and $\theta=30^\circ$. The $H^z$ is applied perpendicular to the skyrmion plane, which is given by $H^z_x \sin \omega t$ with $H^z_x = 0.05 H_z = 0.00285$. The simulated trajectories and velocities $v=(v_x, v_y)$ are shown in Fig. 4(f) and (g), respectively. We find that the trajectories are again oriented almost in the negative $x$ direction; however, for the ferromagnetic resonance mode with $\omega=0.0664$ the trajectory is slanted toward the positive $y$ direction, which contrasts with the case of the $E^z$-active mode. Noticeably, the higher-frequency mode with $\omega=0.0664$ has a much faster propagation of the skyrmion than the other two modes. However, the usage of this mode is not energetically efficient because this mode is not an eigenmode of the isolated skyrmion but a very intense resonance of the vast ferromagnetic magnetizations, which necessarily leads to large energy consumption and considerable rise of device temperatures.

In summary, we have theoretically found that resonance modes of an isolated skyrmion in a ferromagnet can...
be activated by application of AC electric fields through oscillatory variation of the interfacial DMI. The advantage of this method compared with conventional methods using AC magnetic fields is that we can selectively excite skyrmions without activating gigantic ferromagnetic resonances; this results in a significant suppression of both energy consumption and temperature rise. Our result revealed that among the three $E^\omega$-active modes, the mode with combined clockwise and breathing oscillations induces much faster skyrmion propagation than the previously studied breathing mode. Our findings will pave a new way toward the efficient manipulation of isolated skyrmions and thus will be useful for future skyrmion-based devices.

This work was supported by JSPS KAKENHI (Grant No. 17H02924), Waseda University Grant for Special Research Projects (Project Nos. 2017S-101, 2018K-257), and JST PRESTO (Grant No. JPMJPR132A).

1A. N. Bogdanov and D.A. Yablonskii, Sov. Phys. JETP 68, 101 (1989).
2A. Bogdanov and A. Hubert, J. Mag. Mag. Mat. 138, 255 (1994).
3S. Mühlbauer, B. Binz, F. Jonietz, C. Pfleiderer, A. Rosch, A. Neubauer, R. Georgii, and P. Böni, Science 323, 915 (2009).
4X. Z. Yu, Y. Onose, N. Kanazawa, J. H. Park, J. H. Han, Y. Matsui, N. Nagaosa, and Y. Tokura, Nature 465, 901 (2010).
5N. Nagaosa and Y. Tokura, Nat. Nanotech. 8, 899 (2013).
6S. Seki and M. Mochizuki, “Skyrmions in Magnetic Materials” (Springer Briefs in Physics).
7M. Mochizuki, Phys. Rev. Lett. 108, 017601 (2012).
8O. Petrova and O. Tchernyshyov, Phys. Rev. B 84, 241433 (2011).
9T. Schwarzau, J. Waizner, M. Garst, A. Bauer, I. Stasinopoulos, H. Berger, C. Pfleiderer, and D. Grundler, Nat. Mater. 14, 478 (2015).
10M. Garst, J. Waizner, and D. Grundler, J. Phys. D: Appl. Phys. 50, 293002 (2017).
11M. Mochizuki and S. Seki, J. Phys.: Cond. Matt. 27, 503001 (2015).
12M. Mochizuki and S. Seki, Phys. Rev. B 87, 134403 (2013).
13Y. Okamura, F. Kagawa, M. Mochizuki, M. Kubota, S. Seki, S. Ishiwata, M. Kawasaki, Y. Onose, Y. Tokura, Nat. Commun. 4, 2391 (2013).
14M. Mochizuki, Phys. Rev. Lett. 114, 197203 (2015).
15Y. Okamura, F. Kagawa, S. Seki, M. Kubota, M. Kawasaki, and Y. Tokura, Phys. Rev. Lett. 114, 197202 (2015).
16J. Ohe and Y. Shimada, Appl. Phys. Lett. 103, 242403 (2013).
17Y. Shimada and J. Ohe, Phys. Rev. B 91, 174437 (2015).
18D. Hirobe, Y. Shiomi, Y. Shimada, J. Ohe, and E. Saitoh, J. Appl. Phys. 117 053904 (2015).
19S.-Z. Lin, C. D. Batista, and A. Saxena, Phys. Rev. B 89, 024415 (2014).
20G. Finocchio, F. Büttner, R. Tomasello, M. Carpentier, and M. Klüai, J. Phys. D: Appl. Phys. 49, 423001 (2016).
21A. Pert, V. Cros, and J. Sampaio, Nat. Nanotech. 8, 152 (2013).
22F. Ma, Y. Zhou, H. B. Braun, and W. S. Lew, Nano Lett. 15, 4029 (2015).
23K.-W. Moon, B. S. Chun, W. Kim, and C. Hwang, Phys. Rev. Applied 6, 064027 (2016).
24M. Mruczkiewicz, P. Gruszecki, M. Zelent, and M. Krawczyk, Phys. Rev. B 93, 174429 (2016).
25R. H. Liu, W. L. Lim, and S. Urazhdin, Phys. Rev. Lett. 114, 137201 (2015).
26S. Zhang, J. Wang, Q. Zheng, Q. Zhu, X. Liu, S. Chen, C. Jin, Q. Liu, C. Jia, and D. Xue, New J. Phys. 17, 023061 (2015).
27G. Finocchio, M. Ricci, R. Tomasello, A. Giordano, M. Lunatza, V. Pulaifito, P. Burrascano, B. Azzerboni, and M. Carpentier, Appl. Phys. Lett. 107, 262401 (2015).
28W. Wang, M. Beg, B. Zhang, W. Kuch, and H. Fangohr, Phys. Rev. B 92, 020403(R) (2015).
29P. Bak, and M. H. Jensen, J. Phys. C 13, L881 (1980).
30D. Yi, S. Onoda, N. Nagaosa, and J. H. Han, Phys. Rev. B 80, 054416 (2009).
31K. Nawaoka, S. Miwa, Y. Shiota, N. Mizuochi, and Y. Suzuki, Appl. Phys. Express 8, 063004 (2015).
32S. S. Srinivasa, M. Schott, R. Juge, V. Křižáková, M. Belmeguenai, Y. Roussigné, A. Bernard-Mantel, L. Ranno, S. Pizzini, S.-M. Chérif, A. Stashkevich, S. Auffret, O. Boulle, G. Gaudin, M. Chshiev, C. Baraduc, and H. Béa, arXiv:1804.09955.
33F. Jonietz, S. Mühlbauer, C. Pfleiderer, A. Neubauer, W. Münzer, A. Bauer, T. Adams, R. Georgii, P. Böni, R. A. Duine, K. Everschor, M. Garst, and A. Rosch, Science 330, 1648 (2010).
34X. Z. Yu, N. Kanazawa, W.Z. Zhang, T. Nagai, T. Har, K. Kimoto, Y. Matsui, Y. Onose, and Y. Tokura, Nat. Commun. 3, 988 (2012).
35I. Iwasaki, M. Mochizuki, and N. Nagaosa, Nat. Commun. 4, 1463 (2013).
36I. Iwasaki, M. Mochizuki, and N. Nagaosa, Nat. Nanotech. 8, 742 (2013).
37F. Büttner, C. Moutafis, M. Schneider, B. Krüger, C. M. Günther, J. Geilhufe, C. v. Korff Schmising, J. Mohanty, B. Pfau, S. Schaffert, A. Bisig, M. Foerster, T. Schulz, C. A. F. Vaz, J. H. Franken, H. J. M. Swagten, M. Klüai, and S. Eisebitt, Nat. Phys. 11, 225 (2015).
38S. Woo, K. Litzius, B. Krüger, M.-Y. Im, L. Caretta, K. Richter, M. Mann, A. Krone, R. M. Reeve, M. Weigand, P. Agrawal, I. Lemesh, M.-A. Mawass, P. Fischer, M. Klüai, and G. S. D. Beach, Nat. Mater. 15, 501 (2016).
39G. Yu, P. Upadhyaya, Q. Shao, H. Wu, G. Yin, X. Li, C. He, W. Jiang, X. Han, P. K. Amiri, and K. L. Wang, Nano Lett. 17, 261 (2017).
40W. Jiang, X. Zhang, G. Yu, W. Zhang, X. Wang, M. B. Jungfleisch, J. E. Pearson, X. Cheng, O. Heinonen, K. L. Wang, Y. Zhou, A. Hoffmann, and S. G. E. te Velthuis, Nat. Phys. 13, 162 (2017).