Ultrafast phononic switching of magnetization

A. Stupakiewicz, C. S. Davies, K. Szerenos, D. Afanasiev, K. S. Rabinovich, A. V. Boris, A. Caviglia, A. V. Kimel and A. Kirilyuk

Identifying efficient pathways to control and modify the order parameter of a macroscopic phase in materials is an important ongoing challenge. One way to do this is via the excitation of a high-frequency mode that couples to the order, and this is the ultimate goal of the field of ultrafast phase transitions. This is an especially interesting research direction in magnetism, where the coupling between spin and lattice excitations is required for magnetization reversal. However, previous attempts have not demonstrated switching between magnetic states via resonant pumping of phonon modes. Here we show how an ultrafast resonant excitation of the longitudinal optical phonon modes in magnetic garnet films switches magnetization into a peculiar quadrupolar magnetic domain pattern, revealing the magneto-elastic mechanism of the switching. In contrast, the excitation of strongly absorbing transverse phonon modes results in a thermal demagnetization effect only.

In condensed matter systems, the collective excitation modes (for example, the phonons and magnons) define the energy range that determines all the important and intriguing thermodynamic and macroscopic properties of solids, such as electric, magnetic or crystallographic order, and the superconducting transition temperature. Sufficiently strong resonant excitation of these modes can efficiently induce a profound modification of macroscopic properties, including a permanent change in the order parameter.

Control of the crystal environment arguably represents the most universal mechanism to act on magnetization, as it is present in all materials regardless of their magnetic structure. The emerging field of 'straintronics' aims to study how mechanical strains of the atomic lattice in a solid respond to extrinsic forces. One interesting approach to generate such strains is via the anharmonic interaction of different phonon modes, which transfers an excitation of an infrared (IR)-active mode into a coordinate shift along a coupled Raman-active coordinate. This shift, for the length of the excitation (or for the duration of the phonon lifetime), changes the crystal structure and as such, a new state of the material. Such 'nonlinear phononics' promises a new direction for the ultrafast optical control of solids in their electronic ground state, including, for example, ultrafast steering of metal-insulator phase transitions, or (temporarily) changing superconducting or ferroelectric properties. The most important milestone on the roadmap of nonlinear phononics is an ultrafast strain-induced switching of magnetization.

In our study we used yttrium iron garnet (YIG) thin films on a GGG(001) substrate. In this cubic crystal, strong magnetocrystalline anisotropy renders the cube diagonals the easy directions of magnetization. There are thus eight metastable states for the magnetization vector, as numbered in Fig. 1a. To observe the magnetic domain structures in the sample, we used a magneto-optical Faraday microscope at room temperature (Methods and Extended Data Fig. 1). Details of the methods of detection and identification of different magnetic domains in these garnet films have been discussed before.

To provide resonant excitation of the phonon modes in YIG, we used pulses from Free Electron Lasers for Infrared eXperiments (FELIX) in Nijmegen, the Netherlands. We focused the IR beam on the surface of the YIG film. The pulses of FELIX have been shown to be Fourier-transform limited, with their bandwidth.
experimentally tunable in the range of 0.5–2.0%. We used various pulse trains of different repetition frequencies, from 10-μs-long ‘macropulses’ at a repetition rate of 25 MHz to 10-ns-long bunches of about 10 pulses each and finally to single subpicosecond ‘micropulses’. The results obtained by these grossly different excitations were qualitatively identical.

Before irradiation, we saturated the sample magnetization along the diagonal [111] axis, and observed two types of magnetic phase (states 8 and 4 in Fig. 1a) at zero magnetic field with opposite directions of the polar magnetization components and the same in-plane magnetization components (Fig. 1b). After irradiating the garnet film with 14 μm wavelength light, four large magnetic domains were formed with their magnetization directions along the other diagonals of the structure (Fig. 1c,d). We note that for this particular image, we used a 10-ns-long sequence of FELIX micropulses. This allowed us to keep the optical peak intensity safely below the damage threshold while simultaneously obtaining domains large enough to be stable. A similar effect of the accumulation of the impact of multiple pump pulses was used in ref. 23. We also repeated the experiments using single subpicosecond micropulses of FELIX and obtained very similar switching patterns, although stronger focusing and peak intensities approaching the damage threshold were required. The validity of the multipulse approach is supported by the experiments showing that the only effect of repeating single micropulses (on the timescale of seconds) is a slight growth of the domains, if any. Our micromagnetic simulations that we describe below fully support this conclusion, showing only low-amplitude dynamics appearing after pulse repetition and no difference in the final pattern (Supplementary Video 1). To erase a domain pattern with the switched magnetization and to trigger the initial magnetic state, we applied an external in-plane magnetic field of about 10 mT for a short time.

Aiming to reveal the mechanism responsible for the magnetic switching, we tuned the photon energy to different spectral lines while keeping the pump fluence constant. The normalized switched area (Extended Data Fig. 2) obtained from the images of large magnetic domains is plotted as a function of the pump wavelength within the range 10–30 μm (Fig. 2a). The crystal structure of garnets lattice is described by a 160-atom body-centred-cubic unit cell of the OI10 (Ia3d) space group22, which gives rise to an extended set of IR-active phonon modes in this spectral range23–26. We used synchrotron-based IR ellipsometry to accurately determine the phonon spectra of the garnet film (Methods and Extended Data Figs. 3–5). This method is very sensitive to the properties of thin films because of the oblique incidence of light27, and it yields both the absorption spectrum of IR-active transverse optical (TO) phonons and the loss function spectrum. In the spectral range of interest, the absorption spectrum consists of two sets of bands centred around wavelengths of 17 μm and 27 μm (590 cm−1 and 370 cm−1), conventionally assigned to dipole-active vibrations at the tetrahedrally and octahedrally coordinated Fe3+ and Fe2+ sites, respectively, separated by a bandgap in the phonon density of state characteristic of garnet crystals. The multiple shortwave bands around 17 μm primarily involve Fe3+–O stretching vibrations, whereas those in the long-wavelength range are mainly related to bending and stretching Fe3+–O phonon modes, albeit with some contribution from Fe2+ bending modes27–28.

Correspondingly, the spectral dependence of longitudinal optical (LO) phonon modes as derived from the loss function reveals a pronounced resonant behaviour around wavelengths of 14 μm and 22 μm (730 cm−1 and 445 cm−1). It is clear that the behaviour of the switched area is extremely well correlated with these LO resonances. The somewhat smaller effect in the vicinity of 23 μm can be easily explained by the approximately 25% absorption of IR light by water vapour during our experiments performed in air. In addition, we have verified that the polarization of the pump does not influence the observed switching, which proves the longitudinal character of the excited phonons. Note that we do not observe any switching at the resonances corresponding to the TO phonons, even though the absorption of IR radiation is considerably stronger at the TO resonances (Fig. 2a). Instead, in the range of moderate intensities, we observed the appearance of a randomized multidomain pattern. Such a pattern can be naturally explained by the heating of the sample close to the Curie point with subsequent relaxation. Lower intensities showed no change of magnetic state whatsoever, whereas increasing intensity further resulted in structural damage of the sample (Extended Data Fig. 6). From the pattern of the damage, it is evidently caused by the stresses originating from rapid thermal expansion due to the heating.

**Fig. 2 | Resonant phonon mode excitation of magnetization switching.**

**a** Spectral dependence of the magnetization switching and optical phonon modes defined by the loss function Im(−ε−1) (bottom) and absorption d (top) in a YIG film. The normalized switched area is calculated as the ratio of the switched area to the maximal switched area obtained for 14 μm laser irradiation with a fluence of 1.2 J cm−2 (the laser spot was marked by red dashed circle in Fig. 1b). **b** Dependence of the normalized switched area on fluence for 14 μm pump laser irradiation. The dashed line is a guide to the eye. The inset images show the magnetic domains after switching with irradiation of different pump fluence patterns. The threshold of burning (local damage of the sample) in garnet was detected empirically (red shading).
Nevertheless, the associated strains were not sufficient to lead to magnetic switching.

In contrast to the case of the TO pumping, the excitation at the frequencies of the LO modes did lead to the formation of the four-domain switching pattern. This observation directly confirms that the proposed switching mechanism is a result of resonant pumping of specific excitation—LO phonon modes—and not from mere thermal expansion of the lattice. Stroboscopic pump-probe experiments performed using a different laser source prove that the mid-IR excitation can trigger magnetization dynamics on a time-scale shorter than the usual precessional period (Extended Data Fig. 7). In these experiments, the observed dynamics do not depend on the pump polarization, and show a maximum amplitude of the excited precession at the 14 μm wavelength, thus also proving that the excitation mechanism is the same as in switching experiments.

Figure 2b shows how the normalized switched area increases with an increase of the pump fluence. Note that the threshold appearance of switched long-lived magnetic domains after irradiation is determined by a balance between the stable domain size and the diameter of the laser spot. On the other hand, further increase of the pump fluence above 2 J cm⁻² leads to local damage of the film.²⁸

To demonstrate the equivalency of switching starting from various initial magnetic phases, we then repeated these experiments starting with different initial direction of magnetization, set by a brief application of the external magnetic field. Figure 3 shows the acquired images of the switched magnetic patterns for four initial magnetizations along the image diagonals.

To elucidate the leading role of the LO phonon for the switching, compared with the TO one, we note the basic difference between these phonon modes. The atomic vibrations corresponding to the LO phonon are along the sample normal. Along this direction, the growth-induced anisotropy lifts the inversion symmetry of the O₁₀ group²⁹, while the original symmetry survives in the plane of the film, where the TO mode vibrates. Therefore, the atomic potential is considerably anharmonic, amplifying the nonlinearity of phononic interactions.¹²,¹³,³⁰ Moreover, an electric polarization accompanies
the atomic motion, leading to a much stronger interaction of the LO mode with the environment. This accounted for unexpected strong decay of excitations in semiconductor quantum dots, for example.

As a result of strong coupling, the resonant excitation of the IR-active LO mode leads to the displacement of the equilibrium atomic positions corresponding to a coupled Raman mode. This displacement is equivalent to macroscopic crystallographic strains, with an amplitude far exceeding what can be obtained from thermal expansion (see above the discussion of the heating via TO modes). Moreover, the timescale of the LO- and TO-related excitation is considerably different: strong interaction of the LO phonons with the environment and the nonlinear phononic mechanism of strain assure that the strain pulse generated by the LO mode is shorter than the magnetization precession period, which is required for switching (see the micromagnetic simulation in Supplementary Video 1). This is not the case with TO mode excitation, where the timescale is given by the heat dissipation processes.

To understand the magnetic domain pattern induced by such an interaction, we assume that the induced shift of the potential has, similar to the laser spot, a Gaussian profile. If \( \mathbf{u} \) is the displacement vector, the induced strain is given by

\[
\varepsilon_{ij} = \frac{1}{2} \left( \frac{\partial u_i}{\partial x_j} + \frac{\partial u_j}{\partial x_i} \right),
\]

where \( x_i \) and \( x_j \) are coordinates in the film plane. The complete derivations are shown in the Methods, but we mention here that the strains are assumed to be limited to the film plane (given by \( x-y \) coordinates), so that the magneto-elastic energy can be written as

\[
E_{me} = b_1 \left( \varepsilon_{xx} m_x^2 + \varepsilon_{yy} m_y^2 \right) + 2b_2 \varepsilon_{xy} m_x m_y.
\]

In the limit of an isotropic sample, \( b_1 = b_2 \). Using \( E_{me} \) in the micromagnetic simulations (Methods and Extended Data Fig. 8), we show that a single picosecond-long strain pulse leads to a selective reversal of four different areas in a homogeneously magnetized sample (Fig. 3c). At the centre of the spot, where the excitation has the highest amplitude, the reversal is absent. This can be understood from the first term in equation (2), which is isotropic and thus does not modify the state of the magnetic system. In contrast, the second term has quadrupolar symmetry and is directly responsible for the creation of the peculiar domain pattern. It is thus clear that our simple model fully reproduces the observations, and therefore confirms the switching mechanism. We envisage that, through engineering of material parameters such as the magnetic anisotropy, magneto-elastic coefficients or crystallographic symmetries, alternative domain patterns will be realized with different spatially distributed magnetization states (Extended Data Fig. 9).

In the future, ultrafast modification of the crystal field environment, and thus of magnetocrystalline anisotropy, may become the most universal way to manipulate magnetization. Magneto-elastic interactions are present in all materials and thus can be used everywhere, for example in antiferromagnets. Development of a full understanding of how to control magnetocrystalline anisotropy is also important for applications in devices, for example, because this anisotropy is key to the stability of magnetization in magnetic memories.

**Online content**

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The influence of reflection at the film/substrate interface on the obtained phonon spectra is negligible. Extended Data Fig. 4 illustrates that the 7.3 μm YIG-Co film is opaque in the region of the most pronounced poles and zeroes in the dielectric function of the substrate, and can be analysed as a semi-infinite isotropic medium for simplicity. The film thickness was determined by periodic oscillations in the transparency window of the ellipsometric spectra due to interference of multiple reflections within the film (0.2 to 0.5 eV, not shown). The optical constants of YIG were obtained from independent spectrscopic ellipsometry measurements and also used to process Ψ(ω) and Δ(ω) spectra measured on the YIG film on the substrate at different incident angles. To confirm that the solid lines in Extended Data Fig. 3 correspond to the true dielectric function for isotropic YIG, the data were also analysed through a best-match single-film model calculation procedure, as implemented in the Woolam VWASE2 data acquisition and analysis software28. With the knowledge of the optical constants of YIG and the sample thickness fixed at 7.3 μm, we confined the magnetization to appear near the wavelength-by-wavelength regression procedure to the experimental ellipsometric spectra with the same dielectric function of YIG as shown in Extended Data Fig. 3.

Using the dielectric function of YIG, the absorption coefficient spectrum δ(ω) of Extended Data Fig. 5 (black curve) was obtained. The result is consistent with the absorption spectrum of bismuth/gallium-substituted YIG27. By means of ellipsometry, we were able to measure not only the absorption spectrum of IR-active TO phonons, but also the loss function spectrum (red curve), which reveals strong peaks at 445 cm−1 and 730 cm−1 attributed to LO phonon modes in YIG.

Time-resolved pump–probe demagnetization dynamics measurements using near-IR laser pulses. The intensity of the laser pulses with photon energies in the range of 0.7 to 2.0 eV was 5×1017 photons per pulse. The mid-IR spectral range used in this experiment were generated using frequency doubling in a GaSe crystal (0.35 mm), using the output beams of two commercially available, independently tunable optical parametric amplifiers (OPAs) integrated into a single housing (Light Conversion, TOPAS-Twins). The OPAs were pumped by a commercially available amplified Ti:sapphire laser system (Coherent, Astrella) delivering pulses at a 1 kHz repetition rate with a duration of 100 fs and photon energy of 1.5 eV. The OPAs were seeded by the same white light generated in a sapphire crystal, which ensures separately tunable, but phase-locked, output pulses with photon energies in the range of 0.45 eV to 1 eV. As a result, when these output pulses were mixed in the GaSe crystal, the generated mid-IR pulses were carrier-envelope-phase stable and their energy lay in the range of 65 meV to 250 meV with an average pulse duration of around 200 fs. In the experiments, the mid-IR pulses were focused onto the sample surface to a spot with a diameter of about 150 μm, using an off-axis parabolic mirror. The ensuing coherent spin precession was measured, in a conventional pump–probe scheme, by tracking the polarization rotation, imprinted by the magneto-optical Faraday effect, on co-propagating probe pulses at a photon energy of 1.5 eV.

Micromagnetic simulations. Micromagnetic modelling was performed using the finite-difference time-domain software Object-Oriented Micromagnetic Framework (OOMMF) release 2.0a13. Our sample volume was 500×500×7.5 μm3, and was discretized with cells of size 2.5×2.5×2.5 μm3. The effective magnetic field consists of several terms. We neglect the demagnetizing field as the magnetization is relatively small. We also neglect the transverse field because the exchange length is much larger than the exchange field and the Zeeman field. This substantially stronger uniaxial anisotropy field was inserted to compel the magnetization to stay close to the sample plane, thus allowing us to consider the problem in two dimensions only. As a result, the sample effectively has two easy axes along the [110] and [100] directions. The magneto-elastic field was incorporated using the extension developed by Yuhagi et al.41, and will be discussed later.

Pump-induced demagnetization. The magnetic pattern deduced from the images after irradiation for a wavelength of 17 μm, corresponding to the TO phonon peak, does not show the switching of the large magnetic domains. At the maximum intensity in the centre of the Gaussian profile, a multidomain pattern is observed, with the magnetization exhibiting a faint remanent structure originating at the centre of the image (Extended Data Fig. 6). It thus proves that the switching resulting in the peculiar spatial pattern occurs via a non-thermal mechanism. We should also note that the magnetization switching does not depend on the pump polarization within the whole spectral range of the measurements.

Ultrafast excitation of magnetic precession. To prove that the induced strains can trigger the magnetization dynamics on a timescale shorter than the usual precessional period, we carried out stroboscopic experiments on a different laser
source. Although of insufficient pulse energy to excite the switching, the results clearly show an almost immediate onset of magnetic precession started by the mid-IR laser pulses (Extended Data Fig. 7). In these experiments, the observed dynamics also does not depend on the pump polarization within the whole spectral range of the measurements. Changing the wavelength of the excitation demonstrates the maximum amplitude of the excited precession at a wavelength of 14 µm, thus also proving that the excitation mechanism is the same as in switching experiments.

Theoretical model of strain-induced switching of magnetization. We assume that the anharmonic interaction of the phonon modes results in a shift of the in-plane coordinates proportional to the local intensity of the excitation. Such modification of the local coordinates is proportional to the laser intensity and thus follows the Gaussian profile of the laser spot. Although microscopically of different origin, the macroscopic distribution of the strains will be equivalent to the case of a non-uniformly heated object (for example, a cylinder) with an axially symmetric temperature distribution $T(r)$, which was studied in the past.\(^1\)

The displacement in this case is given by

$$u = \frac{1}{q} \frac{1 + e^{-\rho r}}{1 - e^{-\rho r}} \int_{0}^{r} T(r) dr + (1 - 2\sigma) \frac{1}{\rho} \int_{0}^{\infty} T(r) dr,$$

where $q$ is the thermal expansion coefficient, $\sigma$ is the Poisson ratio and $\rho$ is the radius of the cylinder. By assuming that the effect of the absorbed laser fluence can be treated as a spatially Gaussian 'temperature' that is $T(r) = Ac e^{-\frac{(r-r_0)^2}{2}}$ with amplitude $A$, we obtain the final form of the displacement

$$u = \frac{Aq\rho^2}{2} \int_{0}^{\infty} \left( 1 - e^{-\rho r} \right) + (1 - 2\sigma) \frac{1}{\rho} \int_{0}^{\infty} e^{-\rho r} dr.$$

where $r = \sqrt{(x-x_0)^2 + (y-y_0)^2}$, $x_0 = y_0 = 250$ µm and $a = 50$ µm is the Gaussian standard deviation of the pulse. Therefore the full-width at half-maximum of the pulse is 118 µm. All other strain elements are set to zero. With $R$ tending to infinity (that is, we consider not a cylinder but rather a plane), the second term goes to zero. The constant term before the brackets we denote as $\beta$:

$$u = \beta \frac{1}{\rho} \left( 1 - e^{-\rho r} \right).$$

Going to Cartesian ($x, y$) coordinates and calculating the derivatives gives the spatial distribution of strain

$$\varepsilon_{xx}(x, y) = \frac{\varepsilon_{0x}}{\sigma} \frac{1}{\rho} \left[ 1 - e^{-\rho r} \right],$$

$$\varepsilon_{yy}(x, y) = \frac{\varepsilon_{0y}}{\sigma} \frac{1}{\rho} \left[ 1 - e^{-\rho r} \right],$$

$$\varepsilon_{xy}(x, y) = \frac{\varepsilon_{0xy}}{\sigma} \frac{1}{\rho} \left[ 1 - e^{-\rho r} \right].$$

The magneto-elastic energy is then written as

$$E_{\text{me}} = b_I \left( \varepsilon_{0x} m_x^2 + \varepsilon_{0y} m_y^2 \right) + 2b_2 \varepsilon_{0xy} m_x m_y.$$

The spatial distribution of the first (symmetric) and second (antisymmetric) terms is shown in Extended Data Fig. 8. To obtain the ground magnetic state of the sample, we uniformly tilted the magnetization slightly from a selected easy axis. The Gilbert damping parameter $\alpha$ was set to 0.5, and the magnetization was allowed to relax to its equilibrium state (defined as complete when the maximum $\Delta M/dt$ across all cells became smaller than 0.01 T/s). Dynamic simulations were performed using the final magnetic state obtained from the prior relaxation process. Calculations were performed with $\alpha = 0.2$, and the magnetization was saved every 100 fs for a total duration of 50 ps. The laser-pulse magnetic field $h_{\text{ac}}(t) = h_{\text{ac}}(t) = 8.2$ T, and is split into time- and space-varying functions. In time, the pulse has Gaussian character

$$h_{\text{ac}}(t) = \exp \left( -\frac{(t-t_d)^2}{t_r^2} \right),$$

where $t_d = 2$ ps is the pulse’s delay in time, and $t_r = 1$ ps is the pulse’s full-duration at half-maximum. In space, the pulse is of the form $h_{\text{ac}}(x) = -\frac{4k}{\pi \lambda_{\text{ac}}}$. The magneto-elastic energy density

$$U_{\text{me}}(r) = b_{\text{II}} \sum_j \int M_j^2 d\Omega + b_{\text{III}} \sum_i \sum_j M_i M_j r_{ij}.$$

The indices $i, j = 1, 2, 3$ correspond to the Cartesian coordinates $x, y$ and $z$ respectively. We assume that the sample is isotropic, and so $b_{\text{II}} = b_{\text{III}} = -1$. The magneto-elastic strain terms $\varepsilon_{ij}$ are obtained by considering the sample to be non-uniformly heated in the $x$–$y$ plane by a Gaussian pulse. The resulting strain-induced domain pattern is shown in Fig. 3d. In addition, we provide the simulation result using strain pulses. The dynamics of such a process shown in Supplementary Video 1.

Data availability

All other data that support the plots within this paper and other findings of this study are available from the corresponding author on reasonable request. Source data are provided with this paper.

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Author contributions

A.S. and A.K. conceived the project. A.S. performed the magneto-optical imaging together with C.S.D. D.A., K.S. and A.C. performed time-resolved magnetization dynamics measurements. K.S.R. and A.V.B. performed IR ellipsometry measurements. K.S.D. performed the micromagnetic simulations. A.S., A.K., A.V.B. and A.V.K. jointly discussed the result and wrote the manuscript.

Competing interests

The authors declare no competing interests.

Additional information

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Correspondence and requests for materials should be addressed to A.S. or A.K.

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Extended Data Fig. 1 | Magnetization states in YIG. Magnetic domains pattern with four magnetization states in relation to the crystal symmetry in YIG/GGG(001). The large stripe-like domains are formed that have magnetizations along the [1-11] and the [11-1] axes. Within them, small domains are found that possess magnetization along the [111] and the [1-1-1] axes. Magnetization orientations in the domains and type of the domain structure have been identified following the procedure explained in refs. 18,36.
Extended Data Fig. 2 | Profiles of magnetic domains. Top panel from left to right shows the images of domain structures before the laser excitation (a), after irradiation at pump wavelength 14 μm (b) and difference of these images (c). The laser spot was marked by red dashed circle. The initial domain structure was observed by applying an external in-plane magnetic field 10 mT along the [1-10] axis for one second. The graph shows the spatial profiles of the intensity from the difference image with magnetic contrast after irradiation. The profiles of intensity in the azimuthal directions $\phi_1$, $\phi_2$ and $\phi_3$ are marked on image (c). The red solid line represents the beam profile (half part) of irradiation pulse. The images are 210 x 210 μm$^2$ large.
Extended Data Fig. 3 | Phonon spectra of YIG. Complex dielectric function $\varepsilon(\omega) = \varepsilon_1(\omega) + i\varepsilon_2(\omega)$ (black) and the imaginary part of the dielectric loss functions Im($-1/\varepsilon$) (red) of YIG obtained by direct inversion of ellipsometric spectra. Vertical dashed lines indicate zero crossing of $\varepsilon_1(\omega)$. 
Extended Data Fig. 4 | Phonon spectra of GGG. Imaginary parts of the dielectric function $\varepsilon(\omega)$ (black) and the loss function $-1/\varepsilon(\omega)$ (blue) for Gd$_3$Ga$_5$O$_{12}$. The shaded background shows the infrared transmission spectrum of 7.3 μm thick YIG:Co. The 15% to 20% transparency region around the LO mode at 490 cm$^{-1}$ is marked by the dotted lines in Extended Data Figs. 3 and 5.
Extended Data Fig. 5 | Infrared absorption and loss function. Absorption coefficient (black) and imaginary part of the loss function (red) of YIG determined from ellipsometric spectra.
Extended Data Fig. 6 | Domain structures after demagnetization. The images of magnetic domain structures before the laser excitation (a), after irradiation at pump wavelength 17 μm with fluence 1.2 J·cm⁻². (b) and the image of sample damage after irradiation with fluence 2.5 J·cm⁻². The difference images between before and after irradiation with fluence 0.3 J·cm⁻² (d), 0.7 J·cm⁻² (e) 1.2 J·cm⁻² (f). The laser spot was marked by red dashed circle. The white arrows are marked the crystallographic axes in the YIG film.
Extended Data Fig. 7 | Ultrafast magnetization precession in YIG. Time-resolved magnetization precession induced by the femtosecond pump pulses in YIG film. The out-of-plane component of the magnetization is detected with the help of time-resolved magneto-optical Faraday rotation. (a) The magnetization precession measured using different pump wavelengths in the range from 8 μm to 17 μm. (b) Dependence of the normalized precession amplitude on the pump wavelength (black points) and magnetization switching (blue points) from Fig. 2. The normalized precession amplitude is calculated as the ratio of the precession amplitude to the maximal amplitude obtained for 14 μm pump laser-induced with pump fluence 0.06 J·cm⁻². The red solid lines are fits of damped harmonic oscillations. The frequency of the precession for different pump wavelengths is constantly 4.54 ± 0.25 GHz.
Extended Data Fig. 8 | Spatial distribution of magneto-elastic energy. Magneto-elastic energy distribution created by the excitation of the phonon mode: (a) symmetric part with $b_x = 0$ (the arrows on the axes indicate $x$–$y$ coordinates), and (b) antisymmetric part with $b_y = 0$. 
Extended Data Fig. 9 | Simulation of magnetization switching for different anisotropy parameters. Micromagnetic simulations of switching after the application of the strain pulse with different anisotropy environments. All material parameters are the same as those described in the Methods, except for the following modifications. (a) The easy axes of the cubic magnetocrystalline anisotropy are aligned along [100], [010] and [001]. (b) An additional uniaxial magnetic anisotropy (strength $K_u = -0.25\, kJ\cdot m^{-3}$) is inserted with its easy axis aligned along [100]. (c) Same as in (b) except the strain pulse amplitude is reduced by 20%. The arrows and the color-coding indicate the orientation of the magnetization in the $x-y$ plane after the application of the single strain pulse with duration 1ps.