

Spin-wave resonance in yttrium iron garnet stripe domains

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We study a thin-film yttrium iron garnet sample that exhibits magnetic stripe domains due to a small perpendicular magnetic anisotropy. Using wide-field magneto-optic Kerr effect measurements we reveal the domain pattern evolution as a function of applied field and discuss the role of cubic anisotropy in the domain formation. Rich magnon spectra are observed in the stripe domain states, with a range of excitation conditions providing distinct spectra. The measurements are interpreted using micromagnetic simulations to provide the spatial profiles of each resonance mode. We further simulate domain patterns and resonance spectra accounting for the cubic anisotropy, with good correlation to experiment. This study highlights how noncollinear magnetic domain structures can host complex resonant behavior in a low-damping magnetic material, with potential use in future magnonic applications.

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I. INTRODUCTION

Yttrium iron garnet (YIG) is a widely used material within the field of spintronics [1] and magnonics [2–4] due to its low intrinsic damping (Gilbert damping constant α down to 5×10^{-5} [5]). For magnonic applications, YIG is typically studied in the collinear saturated magnetic state where the homogeneous ground state of magnetization is effective in describing the spin-wave excitations [6–12]. However, the magnetic texture of YIG can exhibit more complex noncollinear states [13–16] which have so far been underexplored relative to the collinear state. Nonuniform magnetic textures such as stripe,

helical, and vortex states have been shown to display rich resonant modes with distinct spatially distributed microspin characteristics [17–21]. Spin waves in noncollinear magnetization textures have been actively studied as a platform for unconventional computing applications, e.g., reservoir computing [19,22–25]. Fabrication of YIG for magnonic devices often requires etching, which may lead to surface damage and a reduction in material quality, including enhancing damping [26]. Other techniques such as laser heating [27] and ion beam irradiation [28] still present the inflexibility of making irreversible changes to the physical properties. On the other hand, magnetic domain structures naturally offer a route to reconfigure spin-wave spectra and propagation without permanently altering material shape or properties [29,30]. For example, domain walls can act as nanochannels to guide spin waves [31–33] or as local spin-wave emitters [34–36].

Micromagnetic domain patterns result from minimizing local energetics, depending on factors including magnetic anisotropy, magnetic dipole interaction, and applied field. High perpendicular magnetic anisotropy (PMA) materials

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often exhibit stripes, labyrinths, and bubble domains [37], whereas materials with cubic anisotropy host Landau and fractal patterns in order to reduce the closure energy [38]. Stripe domains are also often found in thin films with small PMA [39–42], often discussed with the condition $Q = 2K_{\perp}/\mu_0 M_s^2 < 1$, where $K_{\perp}$ is the uniaxial anisotropy constant and M_s the saturation magnetization [38]. These stripes can exist at zero external field when the film thickness is above the critical thickness, $D_{\text{cr}} = 2\pi\sqrt{A/K_{\perp}}$, where A is the exchange stiffness. For thicker films a Landau structure of flux closure domains can form at the film surface in order to reduce the surface energy [38].

Resonant properties of the stripe domain state have been studied in Bi- and Ga-substituted YIG [43–45], among other materials [41,46–50]. These studies typically show multiple resonant modes in domains and domain walls, and characteristic modes of flux closure caps that appear in thicker films [46]. Previous studies of the propagation of magnetostatic waves in YIG stripe domains [51–53] found that the different domain states could offer varying modes of propagation with a limited amount of theoretical and numerical analysis using a microspin model.

In this study, we characterize both static and dynamic properties of 3- μm -thick YIG thin films in a range of collinear and stripe domain states under a small ($<10\text{mT}$) magnetic field. For the dynamic properties, we concern ourselves primarily with the standing wave modes that are set up in the cross section of the stripe domain pattern. We observe how this regime enables mode hybridization due to the confined nature and varying field-frequency trajectories of the various magnon modes. Micromagnetic simulations provide spin-wave spectra that agree well with our observations, allowing us to map out the spatial distribution of the various resonant modes within different regions of the magnetic domain pattern. This study sheds light on the enhanced range of microstate and dynamic properties accessible in garnet films.

II. RESULTS

A. Stripe domain states and field evolution

The YIG crystal studied is a 3- μm thin film grown on a [111]-oriented gadolinium gallium garnet (GGG) substrate using liquid phase epitaxy (LPE) [54,55] as illustrated in Fig. 1(a). The PMA in the sample was determined by comparing in-plane and out-of-plane (OOP) hysteresis loops measured by a vibrating sample magnetometer and was found to be $1400 \pm 200 \text{ Jm}^{-3}$ (providing a Q value of 0.11). This PMA causes the formation of stripe domains as confirmed by magneto-optic Kerr effect (MOKE) images shown in Fig. 1(b). Although the crystal structure of YIG does not allow any intrinsic uniaxial anisotropy, since the magnetostriction constant λ_{111} of YIG is negative, a tensile strain will tend to create an easy axis perpendicular to the plane of the film [54]. The lattice constant of GGG (12.383 Å) being slightly larger than that of YIG (12.376 Å) [56,57] is then advantageous for inducing a small OOP anisotropy, though it has also been suggested that growth-induced anisotropy may be the origin of PMA in thick YIG films when the growth direction is [111] [58,59]. Nevertheless, PMA is necessary for the formation

of stripe domains as previously observed [14,16,52,60]. This PMA can be small in comparison to the shape anisotropy, and in the specific case of the Landau domain structures present in this paper it is typical for the uniaxial anisotropy energy to be small compared to the demagnetizing energy and the film thickness to be large compared to the Bloch wall width parameter, $\sqrt{A/K_{\perp}}$.

We begin by examining the remanent magnetization state after saturating along various in-plane angles from the $\langle 1\bar{1}0 \rangle$ direction, which is determined from the substrate manufacturer specification [see Fig. 1(a)]. As shown in Fig. 1(b), for 0° and 60° (directions referred to as effective easy axes), collinear magnetic stripe domains are observed, whereas for 30° and 90° (directions referred to as effective hard axes), a zigzag pattern emerges. As explained below, the sixfold symmetry of this effective anisotropy in the (111) plane is consistent with the cubic anisotropy of YIG ($\langle 111 \rangle$ easy axis). Stripe realignment also occurs for externally applied fields (H_{ext}) below saturation. As such, we ensure that a sufficiently large field was applied to align the stripe domains prior to measurement.

To further understand the evolution of stripe domains, we used MOKE microscopy to image the magnetization state in an external field applied along the $[1\bar{1}0]$ (effective easy) and $[11\bar{2}]$ (effective hard) directions. The evolution of the stripes when the field is applied along the $[1\bar{1}0]$ axis is shown in Fig. 1(c). In this direction the stripe domains are collinear and aligned with the field. As the field strength increased, a gradual decrease of the stripe periodicity was observed from $2.8 \pm 0.2 \mu\text{m}$ to $2.4 \pm 0.1 \mu\text{m}$ prior to domain disappearance at $\approx 5.5 \text{ mT}$ —see the increase in wave number in the reciprocal mapping of Fig. 1(e). In contrast, the domain evolution for fields applied along the $[11\bar{2}]$ axis displays more significant changes as shown in Fig. 1(d). The stripes originate in the zigzag state, which quickly gives way to the field-aligned stripes at a relatively low field before gradually returning to the zigzag state at 3 mT. After this, the domains straighten along the field axis. The periodicity of the stripe domains increases significantly from $2.8 \pm 0.2 \mu\text{m}$ to $5.2 \pm 0.1 \mu\text{m}$, with one set of stripes widening while the other narrows, before eventually saturating at $\approx 5.5 \text{ mT}$ as displayed more precisely in Fig. 1(f), which shows the reduction of the wave number along the direction perpendicular to the applied field. We can also identify the zigzag state via the broadening in reciprocal space along the x axis showing the finite periodicity in this direction [Fig. 1(g)]. While there is a gradual transition from the field-aligned state to the zigzag state, the transition from the zigzag state to field aligned at greater field is much more abrupt. Our observation is consistent with previous work for (111)-planed YIG [14,52,60], although the zigzag state was not present in zero field. The effective easy and effective hard axis behaviors are representative of coherent rotation and domain propagation magnetization processes, respectively. Thus far we have referred to the $[11\bar{2}]$ and the $[1\bar{1}0]$ directions as the effective hard and effective easy axes, respectively, as in reality the hard and easy axes are the $\langle 100 \rangle$ and $\langle 111 \rangle$ directions, respectively. Further, given the crystalline direction of our sample, all the in-plane directions are equipotential, which we will show later on. However, the magnetization in the sample in the small-field regime is not uniform, making the local magnetic calculations much more complex. In this

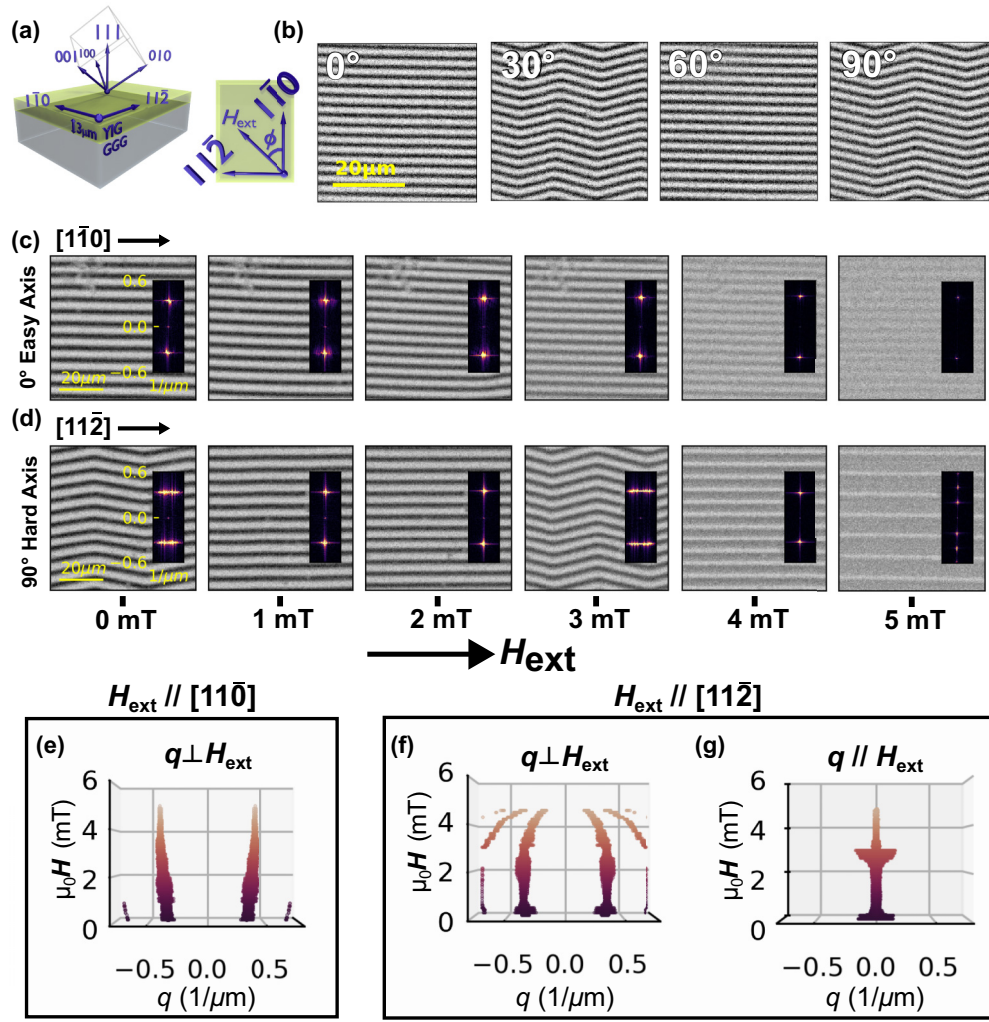


FIG. 1. (a) Diagram of the YIG sample with labeled crystalline axes. (b) Wide-field Kerr images of magnetization state at remanence after saturating $\mathbf{H}_{\text{ext}}$ of 17 mT at 0° , 30° , 60° , and 90° with respect to the $[1\bar{1}0]$ direction. (c), (d) Field evolution of the magnetic state for the field being applied along either $[11\bar{2}]$ direction (c) or $[1\bar{1}0]$ direction (d). Inset shows the Fourier transform of each image. Position of Fourier transform peaks as a function of the $\mathbf{H}_{\text{ext}}$ for the $[1\bar{1}0]$ direction (e) and $[11\bar{2}]$ direction (f), (g). For the $[1\bar{1}0]$ direction, we see a smooth change from stripe to saturated consistent with a coherent rotation magnetization process, whereas for the $[11\bar{2}]$ direction, we see a transition into a state where one set of domains expands while the others shrink, a process which is consistent with a domain propagation magnetization process.

article, we use the terms *effective* hard and easy axes for conciseness in referring to the behavior of these directions rather than in reference to the maxima and minima of the anisotropy energy density.

In order to explain the observed magnetic domain structures, we employ micromagnetic simulations (MUMAX3) [61]. Figure 2(a) shows a vertical section of the thin film with the height of the simulation set to the film thickness and the width to the stripe period. To do this we use a cell size of $16.67 \times 16.67 \times 16.67$ nm with a grid size of $180 \times 180 \times 1$. Periodic boundary conditions were then used to ensure a realistic demagnetizing field. As shown in Fig. 2(b), stripe domains are successfully reproduced as the equilibrium state after removing the field, consistent with our MOKE observation. Flux closure domains are present at the top and bottom surfaces, characteristic of thicker films [38,62]. Results by applying a field along the $[1\bar{1}0]$ and $[11\bar{2}]$ directions agree well with the MOKE images in Fig. 1. In these simulations $\mathbf{H}_{\text{ext}}$ is applied

along the stripe direction (which is perpendicular to the simulated cross section). It should be noted that we only simulate the cross section of the stripe domains. Not only does this limit the wave-vector directions of spin waves that we can observe to a two-dimensional plane, it also prevents the possibility of simulating the zigzag state that occurs when applying a field along the hard axis. As we will show in this article, however, much of the experimental results are captured by these simulations regardless—with larger-scale, computationally more intensive simulations capable of supporting a full long-range zigzag state a fruitful area for future study.

Figure 2(c) shows that the increasing periodicity of the stripes by the in-plane field is reproduced in the $[11\bar{2}]$ simulation. For the $[1\bar{1}0]$ case, we are able to observe a behavior not elucidated by MOKE, which is not sensitive to changes in the magnetization along the thickness direction, i.e., a slanting of the stripe domains under the $\mathbf{H}_{\text{ext}}$ [Fig. 2(d)]. Figures 2(e) and 2(f) show the measured and simulated magnetization loops

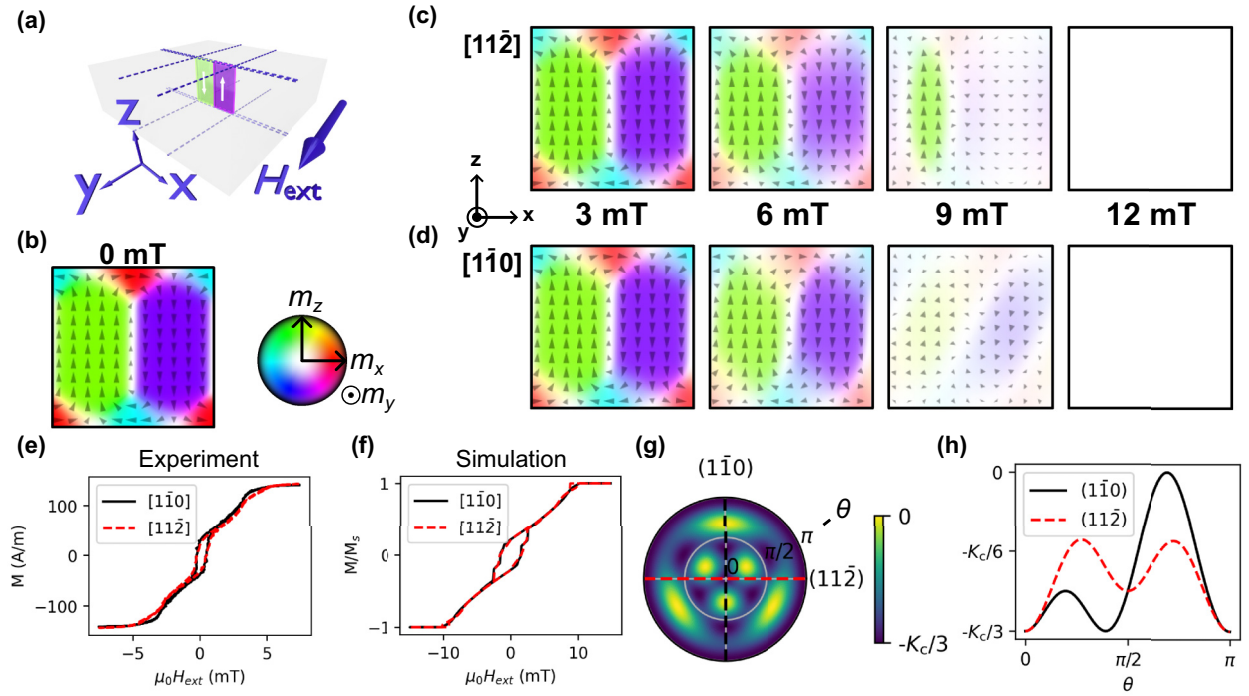


FIG. 2. Simulation of the ground-state magnetization under applied fields. (a) Schematic showing the simulation geometry. (b) Remnant state simulation. With the micromagnetic state for increasing field along the (c) $[11\bar{2}]$ and (d) $[1\bar{1}0]$ axes. (e) Measured and (f) simulated magnetization loops for both the $[1\bar{1}0]$ and $[11\bar{2}]$ axes. (g) Cubic anisotropy energy as a function of polar and azimuthal angle revealing the threefold rotational symmetry of the cubic anisotropy around the OOP axis and the mirror symmetry and lack thereof within the $(11\bar{2})$ and $(1\bar{1}0)$, respectively, shown more explicitly in (h), which displays the polar angle dependence of the cubic anisotropy energy in the $(11\bar{2})$ and $(1\bar{1}0)$ planes.

for the two different axes. While the general shape of these matches well with a small hysteresis in the center of the loop due to the magnetization of the domain walls reversing, one visible disagreement is that the width of hystereses in the simulation is larger than those in experiments, as is the saturation field. We speculate that this is due to the fixed lateral period being imposed on the sample rather than being optimized for every field value, in addition to effects from temperature reducing the coercive field for the experiment compared to simulation—further discussions can be found in Appendix A. Though not having good quantitative agreement, the qualitative agreement is what led us to use these simulation parameters in the present study.

The field evolution of the domain pattern can be explained by examining the first-order cubic anisotropy energy density E [Fig. 2(g)], given by the following [63]:

$$E = K_c \left(\frac{\sin^4 \theta}{4} + \sqrt{2} \cos^2 \phi \cos \theta \sin \phi \sin^3 \theta + \frac{\cos^4 \theta}{3} + \frac{\sqrt{2} \cos \theta \sin^3 \phi \sin^3 \theta}{3} \right), \quad (1)$$

where θ is the polar angle from the $[111]$ direction and ϕ is the azimuthal angle in the (111) plane such that $\phi = 0$ points along the $[1\bar{1}0]$ direction. The three hard axes are directed along the $\langle 100 \rangle$ directions. Using this, we can examine three important cases: the energy density for the (111) , $(1\bar{1}0)$, and $(11\bar{2})$ planes. The anisotropy energy in the (111)

(i.e., the plane that defines the film and contains the $[11\bar{2}]$ and $[1\bar{1}0]$ directions), is an equipotential within the plane expressed as $E_{(111)} = K_c/4$, as we previously mentioned. As such, there is no hard or easy axis within the plane.

We can also calculate the energy density as a function of polar angle from the $[111]$ direction in the $(11\bar{2})$ (containing the $[1\bar{1}0]$ direction) and $(1\bar{1}0)$ (containing the $[11\bar{2}]$ direction) planes, $E_{(11\bar{2})}$ and $E_{(1\bar{1}0)}$, respectively, as follows:

$$E_{(11\bar{2})} = K_c \left(\frac{\cos^4 \theta}{3} + \frac{\sin^4 \theta}{4} \right), \quad (2)$$

$$E_{(1\bar{1}0)} = K_c \left(\frac{\cos^4 \theta}{3} + \frac{\sqrt{2} \cos \theta \sin^3 \theta}{9} + \frac{\sin^4 \theta}{4} \right). \quad (3)$$

From this we can observe that $E_{(1\bar{1}0)}$ has a term that is odd in θ around the in-plane direction, $\theta = \pi/2$ [Fig. 2(h)]. As such, when H_{ext} is applied along the $[11\bar{2}]$ direction (effective hard axis) the two stripes begin to cant toward the film plane. As the in-plane component of the magnetization increases, one stripe domain will approach a maxima of the anisotropy energy while the domain with the opposite OOP component of the magnetization will approach an energy minima. As such, one set of stripes expands while the other shrinks, as one stripe magnetization becomes more favorable than the other. Conversely, when magnetizing along the $[1\bar{1}0]$ (easy axis) direction, we can see that $E_{(11\bar{2})}$ is completely even in θ [Fig. 2(h)], and as such neither stripe is favored. By more closely examining the energy density [Fig. 2(g)], we can also

see that as the stripe is magnetizing into the film plane, it becomes energetically favorable for the magnetization of the stripes to tilt slightly, explaining the behavior close to saturation seen in the simulations.

B. Stripe domain resonances

We performed rf absorption experiments by placing the YIG sample on a broadband coplanar waveguide. The rf transmission coefficient S_{12} was measured as a function of frequency and applied magnetic field using a vector network analyzer. We apply external dc magnetic field either perpendicular or parallel to the oscillating magnetic field, so-called perpendicular and parallel pumping geometry, to capture all the modes excited in the unsaturated regime. We repeat the same experiments after rotating the sample by 90° in-plane. While in simulation the stripes are aligned along the field direction by design, in the experiment we apply an initial saturating field prior to each scan to ensure that the stripes are aligned along the field direction. It should be noted that since the sample is laterally larger than the signal line, an OOP component of the rf excitation for both pumping regimes is present in the YIG sample. Furthermore, because of the large width of the signal line (0.5 mm), we assume a uniform excitation given the comparatively smaller domain width.

The noncollinear nature of the magnetic ground state exhibited by the YIG films leads to a variety of dynamical behavior, producing the rich spectra that we reveal in this paper. The antiparallel ordering of the stripe domains causes the system to exhibit both acoustic and optical modes based on the direction of rf excitation with respect to the applied field [Fig. 3(a)]. Nonuniformity of the domain structure leads to mode localization to regions with aligned magnetization, producing characteristic modes for the bulk domains, flux closure caps, and domain walls [Fig. 3(b)]. The finite thickness of the film allows for the formation of standing-wave modes, where only the modes with a net dynamic component aligned with the uniform rf field are allowed by the excitation, disallowing evenly ordered modes [Fig. 3(c)]. Dipolar and exchange coupling both between and within regions leads to mode coupling, producing an avoided crossing where the modes would cross and creating hybridized modes which exhibit features of both the uncoupled modes [Fig. 3(d)]. This becomes a relatively frequent occurrence due to the variety of dynamic behavior and spatial confinement of the modes.

Figure 4 presents spin-wave spectra and simulation results for the perpendicular pumping with $\mathbf{H}_{\text{ext}}$ applied along the easy axis [Fig. 4(a)]. We divide the spectra in Fig. 4(b) into three distinct field regimes. A low-field regime for 0–3.5 mT (I) has a clear stripe domain mainly polarized along the OOP directions; there is an intermediate field regime from 3.5 to 6 mT (II) which is associated with a slanting of the stripe domains toward the nearest easy axis that only partially points OOP, according to our simulations; finally, the moments are collinear to the applied field in the high-field regime (III) above 6 mT, where the Kittel mode appears. Using the same parameters for the prior static simulations, we calculated the resonant spectra as shown in Fig. 4(c). Three modes labeled as 1–3 in Fig. 4(c) are identifiable in the experimental spectra in a very similar frequency region. Their spatial power

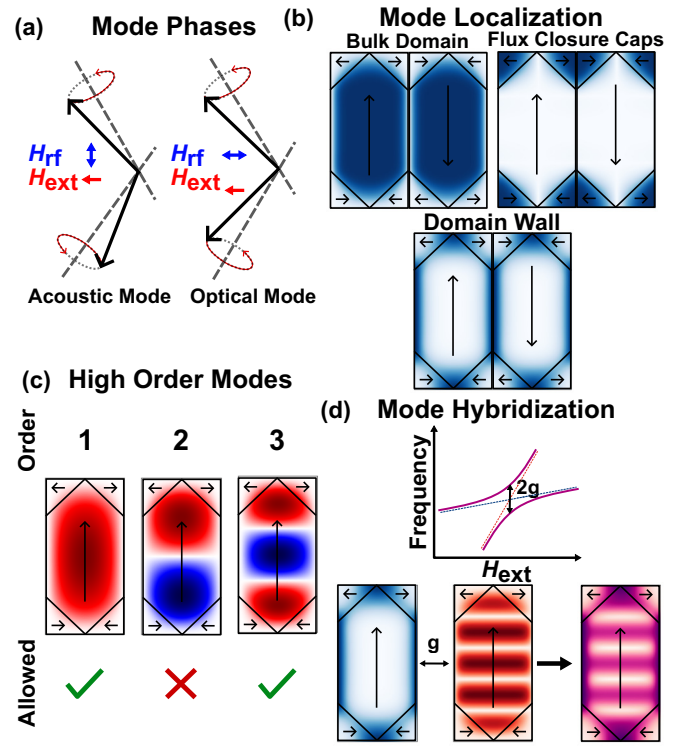


FIG. 3. Schematics of the various dynamic behavior exhibited in the stripe domains. The antiparallel ordering caused by dipole-dipole interactions leads to optical and acoustic modes when the excitation field is parallel or perpendicular to the applied field, respectively (a). Confinement of spin waves within areas of uniform magnetization cause mode localization (b) and the finite-thickness high-order standing-wave modes (c), where the uniform excitation field prevents even-order modes from being excited due to the lack of a net dynamic component. Dipole and exchange coupling between regions of magnetization leads to mode hybridization (d), where two separate modes cross in field and frequency, forming an avoided crossing at this point.

maps [Fig. 4(d)], generated by plotting the Fourier transform amplitude for each cell at a given frequency, reveal that the lowest-frequency mode (Mode 1) is the domain-wall mode associated with the 90° domain walls between the stripe domains and flux closure caps. Mode 2 has strong resonance intensity at the top and bottom surfaces of the bulk stripe domains, which have a tilted magnetization from the bulk moments as determined from the static simulations. Mode 3 is a bulk mode of the stripe domains themselves. The phase maps shown for the bulk mode (3) reveal the two stripe domains to be in phase along the y axis but out of phase along the x axis, which is the axis along which the rf field is applied. This phase relationship is indicative of an acoustic mode, as we would expect for the rf being applied perpendicular to $\mathbf{H}_{\text{ext}}$, and is the case for all the modes excited under this rf direction.

Modes resonating at a higher frequency than Mode 1–3 show a more hybridized character. The modes in Regions 4–6 are much more complex in terms of hybridization, where analyzing individual modes does not offer much insight into the dynamics. Instead, we analyze modes in regions that display some common behavior, which are delineated by a dotted out-

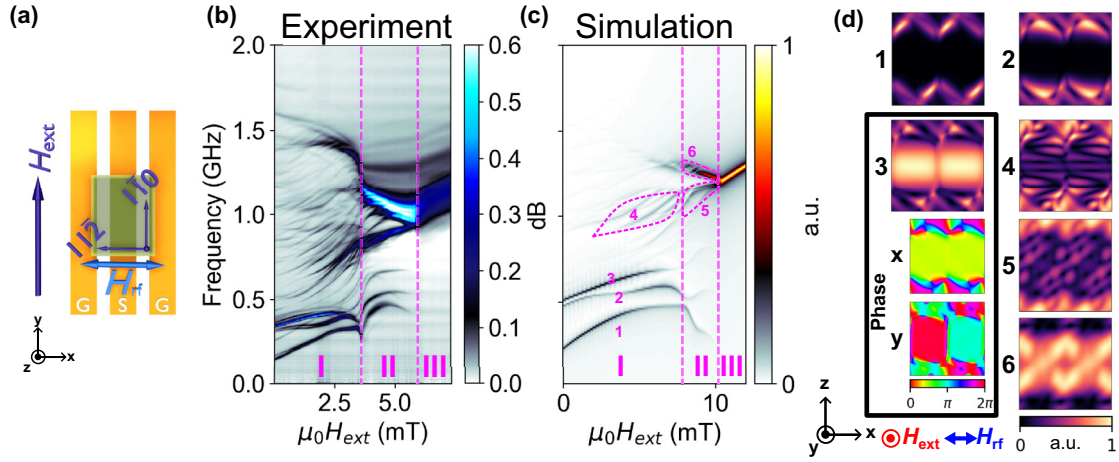


FIG. 4. (a) Diagram showing the relative directions of external field, rf field, and anisotropy axes. S and G denote the ground and signal of the microwave coplanar waveguide. (b) Measured and (c) simulated magnetic resonance spectra of YIG sample for $\text{rf} \perp \mathbf{H}_{\text{ext}}$, which is directed along the $[1\bar{1}0]$ axis. Dotted outlines on (c) delineate regions of hybridized modes sharing some common behavior. (d) Power maps of spin-wave modes labeled on (c) with a phase map shown for the bulk mode (3). The x and z axes are directed horizontally and vertically, respectively, with the y direction being along the axis of the stripe domains (out of the page) and the magnetization within the stripes pointing along the z axis at zero field. These power maps are taken at 3 mT (1–4) and 9.5 mT (5 and 6).

line on the figures. For each region we display a power map of an example mode taken from said region. This hybridization is apparent in the spatial mapping of the example mode for Region 4 in Fig. 4(d), where high-order bulk modes resonate together with modes located in the flux closure domains. The modes in this region are characterized by a finite wave vector parallel to the static magnetization within the stripes [as displayed in Fig. 2(d)], determined through Fourier analysis of the spatial power maps. The two modes from Regions 5 and 6 that meet at the saturation field ($\mathbf{H}_{\text{ext}} \approx 10$ mT) can also be identified in the experiment. The power maps reveal that these modes in these regions are the surface (5) and bulk (6) modes.

Figure 5 shows spin-wave spectra for the parallel pumping configuration and corresponding micromagnetic simulation results for $\mathbf{H}_{\text{ext}}$ along the effective easy axis. Here, we can also identify Modes 1 and 2, which are localized in the same regions as Modes 1 and 2 from the perpendicular pumping results. The difference in field frequency profile is likely explained by the rf field direction, which determines the phase relationship of the dynamic component of the magnetization between the stripes [Fig. 5(d)]. The power maps for these modes also indicates a difference in the hybridizations that occur. Interestingly, these modes primarily associated with the domain walls are not excited efficiently at lower fields

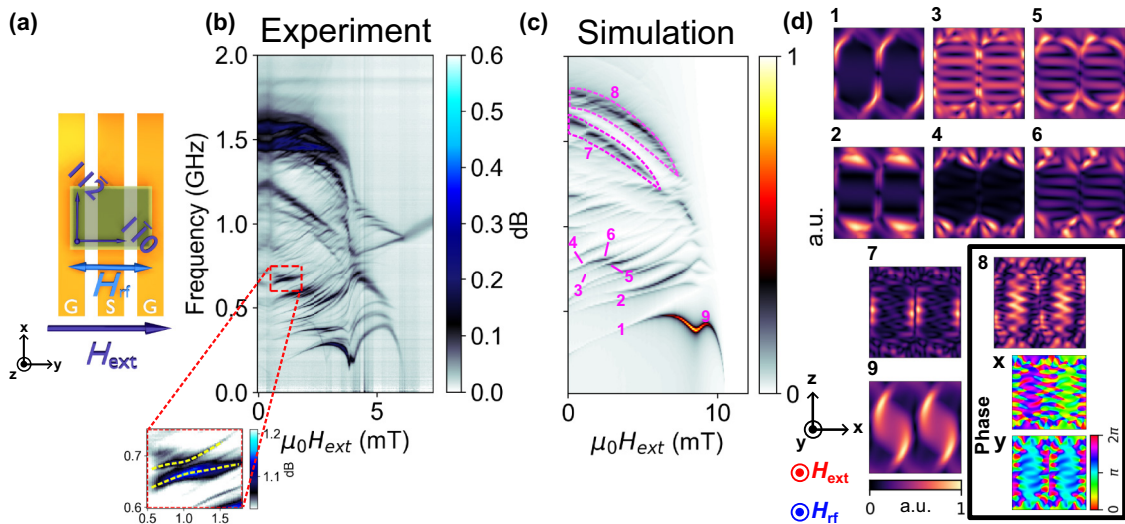


FIG. 5. (a) Diagram showing the relative directions of external field, rf field, and anisotropy axes. S and G denote the ground and signal of the microwave coplanar waveguide. (b) Measured and (c) simulated magnetic resonance spectra of YIG sample for $\text{rf} \parallel \mathbf{H}_{\text{ext}}$, which is directed along the $[1\bar{1}0]$ axis. Dotted outlines on (c) delineate regions of hybridized modes sharing some common behavior. (d) Power maps of spin-wave modes labeled on (c) with a phase map shown for the bulk mode (6). These power maps are taken at 3 mT (1, 2, 5, and 6), 1 mT (3, 4, 7, and 8), and 8 mT (9).

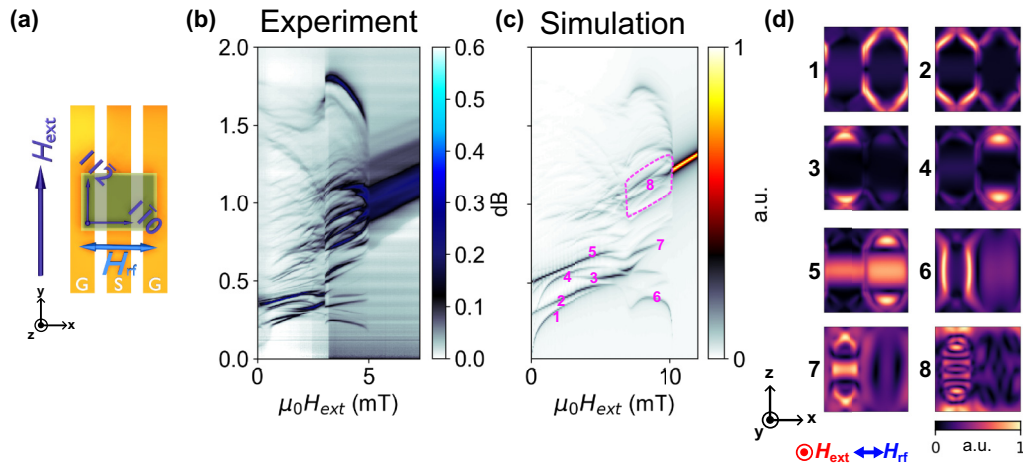


FIG. 6. (a) Diagram showing the relative directions of external field, rf field, and anisotropy axes. S and G denote the ground and signal of the microwave coplanar waveguide. (b) Measured and (c) simulated magnetic resonance spectra of YIG sample for $\text{rf} \perp \mathbf{H}_{\text{ext}}$, which is directed along the $[11\bar{2}]$ axis. Dotted outline on (c) delineates a region of hybridized modes sharing some common behavior. (d) Power maps of spin-wave modes labeled on (c). The power maps were taken at 3 mT (1–5) and 8 mT (6–8).

for parallel pumping. This can be explained by the fact that magnetization in these regions is predominantly parallel to the rf field applied. This aspect does not change as larger fields are applied, suggesting that the hybridization with other modes allows these regions to be more efficiently excited. Modes 3–6 in Fig. 5(d) exemplify the presence of hybridization in our sample. This particular coupling can be observed both in the simulation (as denoted by the numbers 3–6) as well as the experiment (red dotted box). Modes 3 and 4 in Fig. 5(d) show the spatial profiles of the modes away from the anticrossing ($\mathbf{H}_{\text{ext}} = 1$ mT), which reveal these two modes to be an edge mode and the fifth-order thickness modes. For the spatial profile for Modes 5 and 6, either side of the anticrossing ($\mathbf{H}_{\text{ext}} = 3$ mT), we can observe that the two modes have hybridized, both showing aspects of the spatial profile present in Modes 3 and 4. For the higher-frequency modes, we encounter a similar issue as mentioned earlier, whereby the prominent role of hybridization leads us to discuss the modes in the outlined regions (7 and 8) by their common behavior. Modes in Region 7 display resonance mainly in the 180° domain walls, whereas those in Region 8 are associated with bulk modes having a two-dimensional standing-wave nature both parallel and perpendicular to the thickness direction (z coordinate in Fig. 5), coupled to the edge modes. We also show the phase map for Region 8 to reveal an optical mode phase relationship with the precession in phase along $\mathbf{H}_{\text{ext}}$ (y direction) and out of phase along the x direction.

By following Mode 1 to the high-field regime where we see a discontinuous change in the resonant modes, which we label as Mode 9, we can attempt to understand the intermediate phase in our experiment. As we discussed earlier using Fig. 2(d), the stripe domains pick up a magnetization in-plane perpendicular to the applied field, as the cross section of the stripes also slants in the same direction. Evident from the power map of Mode 9 in Fig. 5, the slanting rotates the direction of the k -vector of these spin waves, increasing the wavelength of the first-order mode in tandem with any other higher-order modes.

We now turn to the case where $\mathbf{H}_{\text{ext}}$ is applied along the effective hard axis, which is summarized in Fig. 6, showing a difference in spectra to Fig. 4, particularly in the intermediate-field region. We attribute this to the evolution of static magnetic states, i.e., for the effective easy axis we observed a slanting of the stripes away from the OOP direction as we increased magnetic field, while for the effective hard axis an expansion of one set of stripes alongside the shrinking of the other occurs. This ground-state difference impacts the resonant spectra for the intermediate-field regime (3.5–5 mT in experiment and 7–10 mT in simulation), together with noticeable difference in the low-field regime below 3.5 mT. The most notable feature is the energy splitting of the two stripes for the low-frequency modes [a splitting of Modes 1 and 2 and Modes 3 and 4 in Fig. 6(d)], which share the same relative localization within each stripe. The power maps in Fig. 6(d) reveal Modes 1 and 2 to be domain-wall modes but relating to different stripes magnetization. Mode 2, active in domain walls of stripe magnetized in the positive OOP direction, has slightly higher frequency than Mode 1, which is associated with the stripe magnetized in the negative OOP direction. We see the reverse for Modes 3 and 4, where the domain magnetized negatively along the OOP has a higher energy for this mode localization. The frequency splitting of these modes is explained via the same mechanism that causes differing widths of the two stripes at higher fields [Fig. 2(c)]. The cubic anisotropy creates a symmetry breaking regarding the mirror symmetry with respect to the midplane for the film, as demonstrated in Figs. 2(g) and 2(h), provided that there is some tilt toward the $[11\bar{2}]$ direction. The asymmetry of the cubic anisotropy leads to both a slight difference in the magnetization of each stripe as well as a difference in the anisotropy energy density. Modes 6 and 7 and those in Region 8 reveal how the intermediate state where we observe uneven widths of the stripe domains can affect the resonant spectra primarily by altering the fundamental wavelength of the magnons along the x direction. Mode 6, like Modes 1 and 2, is predominantly excited in the domain wall. However, in

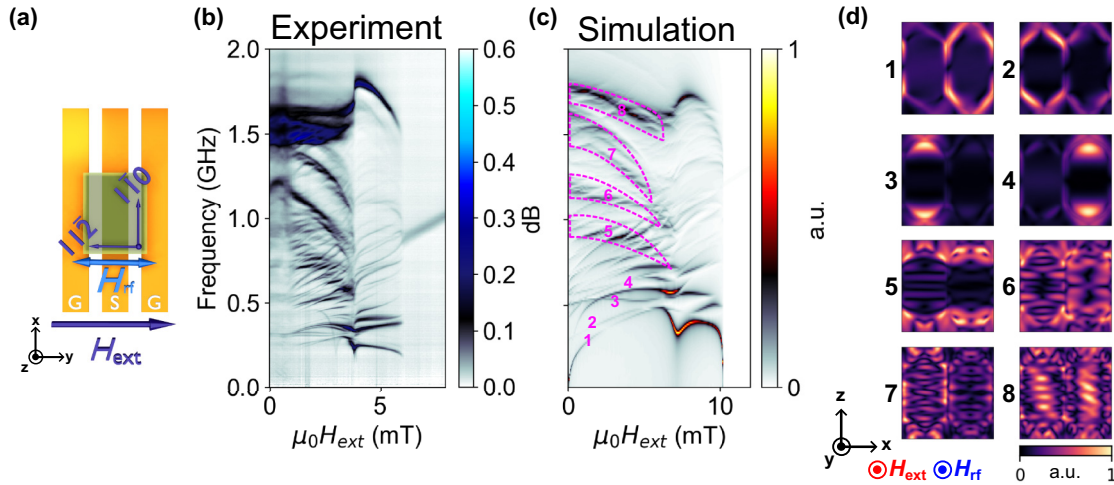


FIG. 7. (a) Diagram showing the relative directions of external field, rf field, and anisotropy axes. S and G denote the ground and signal of the microwave coplanar waveguide. (b) Measured and (c) simulated magnetic resonance spectra for YIG sample for $\text{rf} \parallel \mathbf{H}_{\text{ext}}$, which is directed along the $[11\bar{2}]$ axis. Dotted outlines on (c) delineate regions of hybridized modes sharing some common behavior. (d) Power maps of spin-wave modes labeled on (c). The power maps were taken at 3 mT (1–4), 1 mT (5, 7, and 8), and 2 mT (6).

contrast to Modes 1 and 2, this is localized in the domain-wall region where the magnetization rotation is close to a 180° difference. The complex configuration of dipolar fields in this state makes specific analysis of the modes nontrivial, but certainly the change in domain width in this region alters the wavelength of the standing-wave modes with a k vector directed perpendicular to the domain walls.

Finally, we show the last experimental geometry for the parallel pumping with $\mathbf{H}_{\text{ext}}$ along the hard axis in Fig. 7. The same energy splitting of the lower-frequency modes is present and Regions 5–8 appear to follow the same trend as those in Fig. 5(c). A substantial change in the resonant spectra in the intermediate state is observed as in Fig. 6, though the relative excitation efficiency of positive and negative gradient modes is reversed between the two. This is likely due to the collinearity of magnetization with the external field, as we would expect parallel pumping to be inefficient.

While the general behavior of the dynamics seen in experiment is reproduced by the simulation, there do still remain some discrepancies. A number of possible reasons could be attributed to the source of these differences, such as the choice of anisotropy constants and exchange stiffness used (as detailed in Appendix A), the limitation of the simulation to a single plane, and the periodic boundary conditions imposed. Due to the significant dependence of the resonant spectra on the underlying magnetic ground state, differences between the simulation and experiment in this regard may explain much of the discrepancy in the observed dynamics. It is apparent that for the hard axis case, the simulations do not allow for a changing periodicity seen in experiment [Fig. 1(f)], even though the change in domain size is allowed and observed. Even the more slowly varying periodicity for the easy axis case [Fig. 1(e)] will still have some impact. Other discrepancies in the domain pattern may also exist but are not observable through MOKE, which does not resolve changes across the sample thickness. The tilting of the stripe domains seen in Fig. 2(d) will, in reality, be more or less exaggerated depending on the precise values of K_c and K_L . Nonetheless,

the good qualitative agreement between the simulation and experiment indicates that the most significant aspects of the physical origins are accounted for.

Taken together, the trend observed from the four experiments is that the resonant modes which couple to the excitation in the stripe domain phase are primarily determined by the pumping configuration. Most notably, we selectively excite acoustic and optical modes for perpendicular and parallel rf excitation, respectively. Mode hybridization across the different edge and bulk sections is observed due to local exchange and dipole interactions. The good qualitative agreement of the simulated resonant spectra with the experimental ones demonstrates that, despite the discrepancy of the saturation field, the simulation parameters chosen still allow us to elucidate the origin of the resonant mode observed. At higher fields magnetic anisotropy determines the domain evolution in the ground state, on which different resonant modes are excited. This offers a potential opportunity to enhance the reconfigurability of magnon systems by manipulation of the domain state.

III. CONCLUSION AND OUTLOOK

We have studied the resonant spectra of the stripe domain states in 3- μm -thick LPE-grown YIG thin films. Due to its low damping, we resolved multiple spin-wave modes, where micromagnetic simulations are a powerful tool to understand the spatial distribution of both static and dynamic components of magnetization, e.g., modes populated in the flux closure caps and domain walls. We have shown how the small cubic anisotropy is able to dramatically affect the domain structure as a function of applied magnetic field and therefore the nature of resonant modes on the background magnetization. The PMA present in our sample is not intrinsic to the crystal structure, and as such may be easily controlled by growth conditions, e.g., growth direction in the case of growth-induced anisotropy [58,59]. Application of strain [64,65] or even spin currents [66] leads to potential opportunities for tun-

ing this relatively weak anisotropy, realizing reprogrammable magnonic modes excited on the stripe domain states. Observed resonances depending on the pumping geometry allow for another level of selectivity, controlled by small magnetic fields. This differentiates from other methods such as near-infrared pulses [16]. The stripe direction being influenced by the cubic anisotropy at zero field opens up another pathway to reorient these stripes, and as such to change the excited resonances [67]. Our study is highly relevant to a recent study of magnetic stripes in $\text{La}_{0.67}\text{Sr}_{0.33}\text{MnO}_3$, exhibiting similar resonant and propagation modes tunable via stripe direction [68]. The ability to switch from a stripe domain state to a completely saturated state via the application of a relatively small field creates an opportunity for multifunction devices that make use of both states for different purposes, e.g., data processing and transfer.

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DATA AVAILABILITY

The data that support the findings of this article are not publicly available upon publication because it is not techni-

cally feasible and/or the cost of preparing, depositing, and hosting the data would be prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

APPENDIX A: SIMULATION PARAMETERS

We use MUMAX3 [61] to simulate the experimental results. To do this, we simulate a vertical section of the thin film with the height of the simulation set to the film thickness and the width to the stripe period. Periodic boundary conditions were then used to simulate the demagnetizing field. A schematic of the geometry is shown in Fig. 2(d). For the simulations, we used the anisotropy constants of $K_{\perp} = 1000 \text{ Jm}^{-3}$ and $K_c = -600 \text{ Jm}^{-3}$. The K_c value is an approximate value from literature [69,70], while the $K_{\perp}$ value is chosen to best reproduce the experimental results. This value is somewhat lower than our measured value, and yet the saturation field we see in simulation is approximately double that of the experiment. As the predominant source of anisotropy in our sample is the demagnetizing field. Our method for calculating the uniaxial anisotropy involves subtracting the total anisotropy field from the demagnetizing field that we would expect for a thin film with the given measured value of M_s . Therefore, any error in the value of M_s has a much more significant effect on the calculated K_u . As an example, a 2% change in the M_s value would lead to an over 35% change in K_u .

We also use an exchange stiffness value of $6.5 \times 10^{-12} \text{ J/m}$, a value that is somewhat larger than the commonly reported value of $\approx 4 \times 10^{-12} \text{ J/m}$ [71–73]. While for thin films the stripe domain width is typically equal to the film thickness D , for thicker films the optimal width takes the form of

$$W_{\text{opt}} = 2\sqrt{2D\sqrt{A/K_{\perp}}}. \quad (\text{A1})$$

Using (A1) to calculate the optimal stripe domain width for the lower value gives $W_{\text{opt}} = 1.23 \mu\text{m}$. We then observe that even when the simulation is initiated with stripe domains of period $1.5 \mu\text{m}$ as in the experiment, under applied field these domains will divide to produce a shorter domain width. While a value of $9 \times 10^{-12} \text{ J/m}$ is required for $W_{\text{opt}} = 1.5 \mu\text{m}$, we choose $6.5 \times 10^{-12} \text{ J/m}$ as a balance between these two options such that we do not see the division of the domains under applied field.

APPENDIX B: MATERIAL GROWTH

The YIG film was grown through LPE from a Fe_2O_3 - Y_2O_3 mixture in $\text{PbO-B}_2\text{O}_3$ flux at 910°C . During growth, the GGG substrate was rotated at 100 rpm. To ensure a high-quality film, the solution was stirred at 990°C before LPE growth. To remove flux residue on the YIG surface, the film was cleaned with an etch solution (1:1 ratio of H_2SO_4 and H_2O_2) after the LPE growth.

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