Magnetic hysteresis of dynamic response of one-dimensional magnonic crystals consisting of homogenous and alternating width nanowires observed with broadband ferromagnetic resonance

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We systematically probed the dynamic behavior of homogenous and alternating width (AW) Ni80Fe20 nanowire (NW) arrays using broadband ferromagnetic resonance (FMR) spectroscopy as a function of geometrical parameters such as wire width and interwire spacing. For homogenous width NWs, the FMR responses are markedly sensitive to wire widths and interwire spacing due to spatially varying demagnetizing field. The collective spin-wave mode profile for ferromagnetically and antiferromagnetically ordered ground state has been investigated by controlling the relative alignment of magnetization of neighboring NWs. We show that magnetic ground states of coupled AW NW arrays can be controlled by applying different magnetic field histories, and the collective spin-wave mode is very sensitive to the difference in the widths of wires constituting AW wire arrays. We have also mapped the ferromagnetic and antiferromagnetic ground states magnetic configurations using magnetic force microscopy. Our experimental results are in good agreements with a simple analytical theory we suggest for phenomenological description of the collective oscillations.

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I. INTRODUCTION

The static and dynamic properties of ferromagnetic (FM) nanowires (NWs) of rectangular cross-section have attracted a lot of interest both from a fundamental viewpoint and because of their potential in a wide range of applications such as microwave devices and domain wall logic. They also represent model systems to study the properties of spin-wave excitations in laterally confined magnetic elements. In this regard, both the static and dynamic properties of NW arrays have been studied using various experimental techniques. It has been shown that the magnetization reversal mechanism of NWs is strongly dependent on the geometrical dimensions of the NWs and interwire spacing.

Recently, there has been growing interest in the fundamental understanding of spin-wave (SW) propagation in coupled periodic magnetic NWs called one-dimensional magnonic crystals (MC), which are systems conceived as the magnetic analogue of photonic crystals. Frequency band gaps tuneable by the application of magnetic field have been observed in wirelike structures consisting of shallow grooves etched into a yttrium iron garnet films, one-dimensional array of homogenous NWs separated by an air gap, and synthetic nanostructures composed of periodic arrays of alternating Ni80Fe20 NWs in direct contact with Co NWs.

Recently, we have demonstrated experimentally that a one-dimensional artificial magnetic crystal in the form of a periodical array of magnetic NWs represents an excellent model system to study disorder on periodical lattices. We showed that the degree of disorder of magnetic ground state can be controlled magnetically.

In this paper, we report on a systematic investigation of dynamic responses in tailored homogenous Ni80Fe20 NW arrays as a function of width and interwire spacing using broadband FM resonance (FMR) spectroscopy. We observed that the FMR responses are highly sensitive to the NW width for isolated elements and on magnetostatic interactions for coupled NW arrays. We have also studied the FMR responses of alternating width (AW) NW arrays as a function of differential width of the two NWs constituting the arrays. We show that it is possible to tune the magnetic ground states to obtain both FM and antiferromagnetic (AFM) ordering from the same NW array. The dynamic magnetostatic interaction is evaluated from the frequency shift (Δf0) at remanent state for different ground states. We observed that Δf0 is strongly dependent on the individual wire width in the array. We have also imaged the magnetic configurations of both the FM and AFM ground states using magnetic force microscopy (MFM) and taken representative hysteresis loops using Magneto-Optical Kerr Effect (MOKE) measurement technique.

There is a good agreement between our experimental results and a phenomenological theory we suggest. The theory is based on the model of the Kittel equation. As previously shown for individual wires of infinite length and for arrays of such wires, the initial equations reduce to an eigenvalue problem for an integral or integrodifferential operator (when the exchange interaction is also important). The physical meaning of this result is that the dynamic effective magnetic field, which enters the Landau-Lifshitz magnetic torque equation should have the same shape of the profile across the sample volume as the dynamic magnetization. Hence the proportionality factor between them (or, more generally, a tensor relating them) can be termed the effective demagnetizing factor (or tensor of effective demagnetizing factors). These factors are specific for each particular mode of the standing spin waves across the sample volume (see, e.g., Eq. (10) in Ref. 10 and discussion below that formula.) The obtained effective demagnetizing factors have to be substituted into an equation that follows from the linearized Landau-Lifshitz equation and which can be cast in the form of the Kittel equation. A similar approach has also been used.
for individual prism elements with all three sides comparable in size.\textsuperscript{27} In this latter case, application of an external static magnetic field induces a static demagnetizing field and also results in an inhomogenous profile of static magnetization across the sample. Approximately, this effect can also be cast in the form of an effective demagnetization factor.\textsuperscript{28-30} Furthermore, the approximate dispersion relation for the standing spin-wave modes across the prism [Eq. (13) in Ref. 27] may easily be rewritten in the form of the Kittel equation in which combinations of quantized matrix elements of the dipole-dipole interaction [Eq. (16) in Ref. 27] and of the effective static demagnetization factor along the applied field play the role of respective differences in demagnetizing factors in the Kittel equation.

Extending these known results, in this work we also experimentally demonstrate that the Kittel equation, i.e., parabolic dependence of the square of the resonant frequency on the applied field originally developed by Kittel to describe the magnetic resonance in bulk FM materials also fits quite well experimental data for arrays of dipole-coupled NWs of finite length in the FM collective ground state. We also extend this approach to the case of arrays of antiferromagnetically coupled NWs and demonstrate that the simple model of two antiferromagnetically coupled sublattices produces a line of a shape, which fits well the experimental data in this case. Most importantly, this fit evidences the nature of the collective resonance modes in this case as originating from the hybridization of eigenmodes of two antiferromagnetically coupled sublattices and from the splitting of the respective branches of the dependencies of resonant frequency vs resonant field caused by this hybridization.

The paper is organized as follows. In Sec. II we give details of the samples studied in this work, the method of their fabrication, and the characterization procedures. In Sec. III we show and discuss results of characterization of reference sparse arrays of individual (or isolated) wires. Sec. IV is devoted to the presentation and theoretical discussion of the experimental results obtained on the homogenous width periodic arrays of dipole-coupled nanostripes. A number of experimentally measured dependencies are demonstrated, which provide detailed analysis of collective dynamics in this geometry. We also suggest a simple analytical theory, which is in fair quantitative agreement with the experimental data. The focus of Sec. V is periodic arrays of dipole-coupled alternative width wires. Here we first present the samples that will be experimentally characterized. Then we give the preparation details of the FM and the AFM magnetic ground states for these arrays. Section VA is devoted to the discussion of the dynamic response of the FM ground state and Sec. V B to the AFM state. In the same section we also suggest a simple phenomenological theory of the two antiferromagnetically or ferromagnetically coupled sublattices to explain the experimental results presented in both Secs. V A and V B.

II. DETAILS OF THE EXPERIMENT

Periodic arrays of homogenous and alternating widths Ni\textsubscript{80}Fe\textsubscript{20} NWs of thickness \( t = 30 \) nm and \( L = 10 \) \( \mu \)m were fabricated directly on top of a coplanar waveguide (CPW) using electron beam lithography (EBL) followed by electron-beam deposition and lift-off process. Scanning electron microscope (SEM) was used to verify lift-off completion and dimensions of the structures. Fig. 1(a) shows the SEM image of a typical CPW showing the ground-signal-ground (G-S-G) lines. The CPW fabricated using metallization and lift-off process consists of 300-\( \mu \)m-long, 20-\( \mu \)m-wide, and 150-\( \mu \)m-thick Au signal lines. The Ni\textsubscript{80}Fe\textsubscript{20} NWs arrays of widths \( w \) in a range from 120 nm to 540 nm were fabricated on top of the CPW signal lines. The interwire spacing was also varied from 80 nm (for the isolated NWs). Shown in Figs. 1(g)–1(i) are coupled alternating NWs arrays with (g) \( w_0 = 200 \) nm; \( w_a = 240 \) nm, (h) \( w_0 = 200 \) nm; \( w_a = 380 \) nm, and (i) \( w_0 = 200 \) nm; \( w_a = 540 \) nm.

**FIG. 1.** (a) SEM image of the CPW showing the G-S-G lines (the SEM of 30-nm-thick Ni\textsubscript{80}Fe\textsubscript{20} NWs with width \( w = 120 \) nm and interwire spacing \( s = 180 \) nm is shown as an inset). SEM images of homogenous NWs with (b) \( w = 240 \) nm; \( s = 360 \) nm, (c) \( w = 380 \) nm; \( s = 570 \) nm, (d) \( w = 540 \) nm; \( s = 810 \) nm, (e) \( w = 540 \) nm; \( s = 120 \) nm, and (f) \( w = 540 \) nm; \( s = 80 \) nm. Shown in (g)–(i) are coupled alternating NWs arrays with (g) \( w_0 = 200 \) nm; \( w_a = 240 \) nm, (h) \( w_0 = 200 \) nm; \( w_a = 380 \) nm, and (i) \( w_0 = 200 \) nm; \( w_a = 540 \) nm.
differential widths (\(\Delta w = w_s - w_0\)) are therefore in a range of 
\(-40\) nm to \(340\) nm. To ensure that the NWs are magnetically 
coupled together, the interwire spacing is kept at \(s = 80\) nm 
for the AW NWs.

The FMR responses were measured in the 1–20 GHz range 
using a broadband microwave vector network analyzer (VNA, 
Agilent 8363C). To obtain a high-frequency response, 
the VNA is connected to CPW by a G-S-G-type microwave probe 
similar to that reported in Ref. 5. The magnetic radio-frequency 
 rf field \(h_f\) of the CPW is applied along the \(Y\) direction 
(hard-axis of the NWs), while an external static magnetic 
field (\(H_{\text{app}}\)) is along the \(X\) direction (easy-axis of the NWs) 
as shown in Fig. 1(a). FMR measurements were performed at 
room temperature by sweeping the frequency for fixed \(H_{\text{app}}\). 
This transmission line measurement was repeated for a number 
of different \(H_{\text{app}}\)-values starting from negative saturation field 
\(-H_{\text{sat}} = -500\) Oe, passing through zero, and then gradually 
increasing to a maximal field \(H_{\text{max}} \leq H_{\text{sat}}\) (“forward half of 
a loop”). The field is then subsequently decreased to \(-H_{\text{sat}}\) 
(“backward half of a loop”). This magnetization protocol can 
be considered as a minor hysteresis loop when \(H_{\text{max}}\) is smaller 
than the coercive field, while it is a full-loop measurement 
when \(H_{\text{max}} \geq H_{\text{sat}}\). The magnetic spin configurations of the 
NWs were imaged using the MFM imaging in the phase 
detection mode with a commercial CoCr coated Si cantilever 
tips magnetized along the tip axis. The scan height was 
maintained at a constant distance of 70 nm.

III. ISOLATED NANOWIRES

In this section we show the results of measurements on 
sparse arrays of permalloy NWs. The interelement separation 
fixed at \(s = 1.5 \times w\) to reduce the effects of magnetostatic 
interactions by dynamic stray magnetic fields.\(^{12,21}\) As follows 
from Fig. 6 in Ref. 12, the residual dipole interaction should 
be insignificant in this case. On the other hand, as one will 
see from Sec. V, this distance is close or equal to the distance 
between the elements of the same sublattice on the arrays of 
AW wires. In this way these arrays model the weak dipole 
interaction within sublattices and thus provide good reference 
data for the phenomenological model we suggest in Sec. V.

We stress that the magnetic dynamics of individual (or 
practically uncoupled) wires has been studied extensively in 
previous years.\(^{25,32}\) So, the main goal of this section is to pro-
vide the reference data for the experimental study of the per-
iodic arrays of strongly dipole-coupled wires (Secs. IV and V). 
However, for completeness of the study, we also provide some 
analysis of these data, which may be useful.

We studied the responses from eight different NWs of width 
in the range from 120 nm to \(540\) nm. Shown in Fig. 2(a) are the 
FMR absorption traces of NWs with \(w = 120, 240,\) and 540 nm 
taken at remanence (\(H_{\text{app}} = 0\) Oe) after first saturating the NWs 
along the easy axis. The corresponding SEM micrographs of the 
NWs are shown in Figs. 1(b)–1(d). As one sees from 
Fig. 2(a), the FMR response is characterized by domination 
of one peak, which is easily identified as the fundamental 
(nodeless) mode of standing spin waves across the wire width. 
The frequency of the fundamental mode at remanence \(f_0\) is 
plotted in Fig. 2(b) for various \(w\)-values. In full agreement 
with the previous studies, the resonance frequency decreases with 
increase in the wire width.\(^{25}\) It is dependent on the NW width 
due to the spatially varying dynamic magnetization, an effect 
often referred as the effective dipole pinning of magnetization 
at the wire edges.\(^{25}\)

Also shown as insets in Fig. 2(b) are the resonance 
frequencies as functions of applied field for \(w = 120, 240,\) 
and 540 nm. For clarity, only the data taken for the lower 
halves of the respective hysteresis loops (i.e., from negative
to positive saturation) are shown. The hysteresis phenomenon for the FMR frequency was observed for all wires structures with a sharp frequency jump at around $H_{app} = H_s$ as indicated in the figure.

As stated in the introduction, the shape of the dependence of the resonance frequency on the applied field should follow Kittel’s parabolic law for the square of the frequency:

$$\omega_0 = \gamma \left( [H_k + (N_{xx} - N_{zz})4\pi M_s] \times [H_k + (N_{yy} - N_{zz})4\pi M_s] \right)^{1/2},$$

(1)

where $\gamma$ is gyromagnetic ratio of the material, with typical values of $\gamma$ for Ni$_{80}$Fe$_{20}$ being $\gamma = 2\pi \times 2.93$ GHz/kOe, and $4\pi M_s$ is the magnetization of the sample along $H_{app}$ direction. It can be replaced with the saturation magnetization $4\pi M_s = 10546$ Oe in saturated state, which was extracted using the least-squares-fit method from FMR data obtained on a reference continuous film. The quantities $N_{xx}$, $N_{yy}$, and $N_{zz}$ are effective demagnetizing factors along the NW length ($L$), width ($w$), and thickness ($t$), respectively. As shown in Refs. 10 and 27, the effective demagnetizing factors vary for different resonant modes, which represent standing spin waves across the volume of the wire. The homogenous driving microwave field of FM resonance experiment is able to detect only the modes with a significant net dynamic magnetic moment. In our FMR experiment, we reliably detect only the fundamental (quasiuniform precession) mode, which possesses the largest net dynamic moment. Therefore, below we will be interested only in the modes that represent quasiuniform precession inside individual wires or (in the next sections) collective oscillations formed on the basis of the quasiuniform precession in individual wires. For this reason below we will term the effective demagnetization factors for the quasiuniform precession mode just demagnetizing factors.

We proceed in the following way. We extract the demagnetizing factors from the experimental data using the least-squares-fit procedure. These data will be used as reference in the next sections. Because various simple analytical models based on approximate values for demagnetizing factors have been popular in recent years to explain experiments on dynamical effects in the nano-objects of prismatic shape, for completeness in this study, we also made calculations using the approximate theory for the isolated wires, all three calculations confirmed that $N_{xx}$ should increase with increase in $w$. We checked this idea with calculations using the approximate theory for the isolated wires, for dipole-coupled wires for $s = 1.5 \times w$, and with an simulation of the ground state for isolated wires for the complete magnetic saturation by the LLG Micromagnetics Simulator. (In the latter case we defined $N_{xx}$ as a ratio of the mean value of the static demagnetizing field $H_s$ across the element volume to the mean value of the $x$-component of static magnetization, which is in accordance with the idea in Refs. 28–30.) All three calculations confirmed that $N_{xx}$ should increase with an increase in $w$ [Fig. 2(c)].

The demonstrated invalidity of Eq. (3) in our case may be explained in the following way. Examining the derivation of Kittel’s equation one finds that the $N_{yy}$ and $N_{zz}$ components enter it as a relation between the dynamic magnetization and the dipole field it produces. Theoretically they are obtained as a solution of a linear eigenvalue problem for the dipole operator (see Refs. 10 and 25). The demagnetizing factor along the applied field $N_{xx}$ has however a completely different origin. In theory it originates from the solution of the separate nonlinear equation describing the static equilibrium. Thus, for FMR on nonellipsoidal bodies of finite length ($N_{xx} \neq 0$), the Eq. (3) may not stand in general.

All the theoretical models we use demonstrate that except for the widest wires $N_{xx}$ is significantly smaller than both $\Delta N_{yx}$ and $\Delta N_{zx}$ [Fig. 2(c)]. This allows one to qualitatively analyze the extracted data assuming $N_{yy} \equiv \Delta N_{yx}$ and $N_{zz} \equiv \Delta N_{zx}$. Under this assumption we see that as the NW widths increases the effective dynamic demagnetizing factors $N_{yy}$ diminishes and the out-of-plane one ($N_{zz}$) grows. The decrease in $N_{yy}$ is in agreement with a drop in the wavelength of the standing spin wave across the wire width (i.e., along $Y$) with an increase in the width of the wires. The decrease in the wavelength decreases the in-plane dipole field of the standing wave (given by $N_{yy}$) [Eq. (17) in Ref. 27] and increases its out-of-plane field (given by $N_{zz}$).

Finally, also for the completeness of this study in Fig. 2(c), we plot $\Delta N_{xy}$ and $\Delta N_{zx}$ calculated using the approximate theories from Refs. 29 and 30. One sees that there is a significant disagreement between all these theories and the experimental data. The main reason for this disagreement

Important, an attempt was made to complete definition of the factors using the condition:

$$N_{xx} + N_{yy} + N_{zz} = 1$$

(3)

and in this way to extract absolute values for all three diagonal components of the tensor of the effective demagnetizing factors. This condition is valid for demagnetizing factors of ellipsoidal bodies. It also stands for the stripes of infinite length magnetized along their length, where the effective demagnetizing factors all are of dynamic origin and for the static demagnetizing factors for the prism geometry.

$$N_{xx} + N_{yy} + N_{zz} = 1$$

(3)
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investigate this effect within the hysteresis loops for these materials. We take FMR data for a number of NW arrays consisting of the wires of the same width \( w = 540 \) nm. The interwire spacing \( s \) is varied from 80 nm to 810 nm. Shown in Figs. 1(d)–1(f) are the respective SEM images for the arrays with \( s = 810, 120, \) and 80 nm, respectively.

Shown in Figs. 3(a)–3(c) are the full hysteresis loop FMR absorption spectra for these three arrays. These data correspond to the forward half of a full loop (i.e., the applied field was gradually increased from \(-500 \) Oe to \(+500 \) Oe). As seen from these panels, in the saturated state the FMR response is dominated by a single peak, while multiple absorption peaks are visible for \( s = 120 \) nm and 80 nm when \( H_{\text{app}} \) is around 100 Oe. Figure 3(d) shows the MOKE result for \( s = 80 \) nm from comparison of Figs. 3(c) and 3(d), we infer that the multiple peaks correspond to the field range of magnetization switching between the two saturated states of the NW array.

Shown in Fig. 3(e) are the representative absorption curves taken at \( H_{\text{app}} = 0 \) as a function of \( s \). Similar to the previously studied well-saturated state, the resonance frequency decreases as the \( s \) is decreased due to the effects of magnetostatic interaction between the wires.

We use the theory of demagnetizing factors for periodic arrays of dipole-coupled nonellipsoidal elements by Tartakovskaya to compare our experimental data with the theory. The black dots in Figs. 3(b)–3(c) are the result of the calculation of resonance frequencies using Eq. (1) and the demagnetizing factors obtained from this theory. One sees a good qualitative agreement in terms of trend with the experimental results. Importantly, there is a noticeable quantitative discrepancy between the theory and the experiment: the calculated frequencies are always slightly higher than the measured ones. A reasonable explanation for the discrepancy is the effective dipole pinning of magnetization at the wire edges and the respective modification in the effective demagnetizing factors, which this simple theory does not account for (see discussion in Sec. III). Note that the interwire dipole coupling decreases the effective dipole pinning of magnetization for FMR (compare the thick solid and the thick dashed lines, which are for coupled and isolated wires, respectively, in Fig. 5 in Ref. 12). Thus, the accuracy of the approach by Tartakovskaya should improve with decrease in the wire separation \( s \).

Inside the major hysteresis loop, we observed multiple absorption peaks as shown in Figs. 3(b) and 3(c). The field range \((\Delta H_{t})\) over which the multiple absorption peaks are observed is found to be sensitive to \( s \). For \( s = 80 \) nm, \( \Delta H_{t} = 90 \) Oe, and for \( s = 120 \) nm, \( \Delta H_{t} = 55 \) Oe. This transition region has also been observed for coupled wires with a bistable magnetization state, for an opposing \( H_{\text{app}} \). As the applied field \( H_{\text{app}} \) is increased, the Zeeman energy increases because the NWs magnetization is antiparallel to the applied field. Some of the NWs will switch earlier than the others because of a smaller energy barrier for magnetization switching due to the presence of defects and irregularities. The dipolar interaction is negligible in the case of isolated NWs; therefore, the magnetization reversal for each wire is determined by its own reversal energy barrier, which is similar for all the NWs. Thus the \( \Delta H_{t} \) is almost 0 as shown in Fig. 3(a). In coupled NWs array, however, the antiparallel alignment of magnetization in

![FIG. 3. (Color online) 2D absorption spectra of homogenous width NW arrays. Wire width is 540 nm. Wire separations are (a) \( s = 810 \) nm, (b) \( s = 120 \) nm, and (c) \( s = 80 \) nm. The MOKE results for \( s = 80 \) nm are shown in (d). The representative FMR spectra for this geometry are shown for \( H_{\text{app}} = 0 \) Oe in (e) and for \( H_{\text{app}} = 120 \) Oe in (f). Filled squares: calculation with the values of effective demagnetizing factors obtained using the theory in Ref. 30.](image-url)
nearest NWs is energetically preferable due to dipole coupling. This stabilizes the antiparallel alignment ("anti-ferromagnetic order"). Higher $H_{app}$ values are required to reverse the whole NWs array of dipole-coupled wires, resulting in an increased $\Delta H_s$ when the $s$ is decreased as shown in Figs. 3(b)–3(c).

As recently shown, the mode with lower frequency originates from compact clusters of wires consisting of unswitched NWs, while the higher frequency mode is the response of clusters of switched NWs. Shown in Fig. 3(i) is a set of the absorption spectra for three arrays of 540-nm-wide NW with different wire separations. The applied field is $H_{app} = 120$ Oe, which corresponds to nearly the middle of the transition regime between the two final FM ground states where multiple modes are observed for the coupled NW arrays. One sees that the intensity of the lower-frequency mode decreases with increase in $s$. This implies that with increase in $s$ the stabilization of the antiparallel (AFM) state becomes less probable because of the decrease in the static dipole interaction between wires. Based on the results of Ref. 23, we may suggest that the number and the sizes of clusters with FM magnetic order decrease with increase in $s$.

We also performed extensive investigation of the effects of magnetostatic coupling by fixing interwire spacing at $s = 80$ nm, while varying the width of the NW in in the range from $w = 160$ nm to 540 nm. For all the coupled NWs, we observed similar multiple modes at the corresponding transition fields. We found that $\Delta H_s$ decreases as the $w$ is decreased.

V. ALTERNATING WIDTH NANOWIRE ARRAYS

A systematic investigation of the FMR responses of AW NW array is presented in this section. Six different types of arrays consisting of wires of two different widths alternated on the structure period have been fabricated. One wire width was fixed at $w_o = 200$ nm, while the width of second type of wires $w_s$ was varied from 160 nm to 540 nm. In all the cases, the interwire spacing was fixed at $s = 80$ nm to allow for maximum coupling. The differential width $\Delta w = w_s - w_o$ varied from $-40$ nm to $340$ nm. Shown in Figs. 1(g)–1(i) are representative SEM images for alternating NWs with $w_s = 240$ nm ($\Delta w = 40$ nm), 380 nm ($\Delta w = 180$ nm), and 540 nm ($\Delta w = 340$ nm).

Shown in Figs. 4(a)–4(c) are representative FMR absorption spectra from a set of full loop FMR measurements for these arrays. It can be clearly seen that the FMR responses are strongly dependent on $\Delta w$. Starting from the negative saturation, all the NWs are aligned along the field direction. A stepwise jump upwards in frequency of the fundamental mode is observed when the field magnitude is decreased. This constitutes an important difference in behavior with respect to the homogenous NW arrays [$\Delta w = 0$; Figs. 3(a)–3(c)] for which the curvature of the field dependence but not the frequency itself changes discontinuously.

For the AW arrays it has been established that the NWs of larger width switch earlier than the narrower ones when the sample is released from the magnetic saturation. The stepwise jump in the frequency is attributed to the realignment of the vector of static magnetization. Shown in Fig. 4(d) are the sketches of different magnetization states of AW arrays corresponding to the field range shown in Fig. 4(a). In the region marked “I” all the wires are magnetized in the $-Y$ direction. This magnetization arrangement is called FM order or FM ground state. The AFM ground state occurs in Region II, in which magnetization vectors in the wider NWs and narrow NWs are antiparallel to each other. In Region III, magnetization in the narrow wires has been already switched, which restores the FM order. The new FM ground state is characterized by orientation of static magnetization in both wide and narrow wires in the $+Y$ direction. The corresponding hysteresis loops shown in the lower panel of Figs. 4(a)–4(c) are in good agreement with the FMR results. The magnetic hysteresis (M-H) loops are characterized by a two-step process because of switching of magnetization in sublattices of wide

FIG. 4. (Color online) 2D absorption spectra and MOKE results for AW NW arrays with different differences in width between the wide and narrow wires: (a) $\Delta w = 340$ nm, (b) $\Delta w = 180$ nm, and (c) $\Delta w = 40$ nm (the MFM image for the FM ground state is shown in the inset). Shown in (d) are the sketches of different magnetization states for NWs corresponding to the field range shown in (a). The MFM image of the AFM ground state is shown in (e).
and narrow wires separately and at different applied fields. A drop in the net magnetic moment occurs around 100 Oe (or 150 Oe or 200 Oe) for \( \Delta w = 340 \) nm (or 180 nm or 40 nm, respectively) when \( H_{\text{app}} \) is increased. The magnetization drop is proportional to the volume fraction of wires, which reverse their magnetization direction.

In order to directly observe the FM and AFM magnetic ground states, MFM images were taken for NWs at remanence for various magnetic field histories. Shown as an inset to Fig. 4(c) is the MFM image taken at remanence after saturating the NWs in the field \( H_{\text{app}} = -500 \) Oe. At the ends of the wires (near the top and the bottom of the image), one sees horizontal lines of the brighter and darker contrast with respect to the average brightness of the image (“end lines”). The uniform contrast along these lines suggests that magnetization in all the wires is oriented in the same direction. Figure 4(e) shows an MFM image taken for a magnetic field history known as suitable to achieve an AFM magnetic ground state. After first saturating the NWs in a field of \( -500 \) Oe, the \( H_{\text{app}} \) was then increased to \( H_{\text{max}} = 227 \) Oe. After this, the field was removed \( (H_{\text{app}} = 0) \) and an MFM image was taken. The alternating contrast along the end lines clearly shows that the magnetization vectors in the neighboring wires are aligned antiparallel to each other.

### A. FM Ground State

Now, let us focus on the FM regions of the absorption spectra shown in Figs. 4(a)–4(c). Interestingly, only one large-amplitude absorption peak can be observed in this field region although there are wires of two different widths on the array. This evidences strong magnetostatic interaction between the wires for small interwire spacing. The collective mode, which can be driven by the spatially uniform microwave magnetic field of the coplanar line is the in-phase or acoustic mode. This mode is characterized by in-phase precession of magnetization in the adjacent NWs. The other fundamental collective mode formed from resonances in uncoupled wires (“individual resonances”)—the optical mode, which is characterized by the antiphase precession of magnetization in the wide and narrow wires—cannot be excited efficiently by the uniform field of the coplanar line and therefore does not show a large-amplitude response.

The field dependence of FMR frequency for the FM ground state is shown in Fig. 5(a) as a function of \( \Delta w \). The data for the reference homogenous \( w = 200 \) nm and \( s = 80 \) nm NW array are also included in the figure. One sees that the resonance frequencies for \( \Delta w = +40 \) nm are closer to the ones for the reference sample than for \( \Delta w = -40 \) nm. We also observed that the frequency of the collective resonance shifts downwards with increasing \( \Delta w \). This may be explained by the downshift in the resonance frequency for the individual resonances in uncoupled wires (Sec. III) with increase in the wire width. The decrease in the frequency of the underlying individual resonances drags the frequency of the collective mode downwards.

We have also compared the resonance frequencies at remanence for arrays of homogenous and AW NWs [Fig. 5(b)]. We observed a frequency difference of 0.15 GHz between the data for AW NW array with \( \Delta w = -40 \) nm and for the homogenous NWs array with \( w = 160 \) nm. However, the fundamental mode for the NWs array with \( w = 200 \) nm is lower in frequency by 0.32 GHz than the one for the AW array with \( \Delta w = -40 \) nm. This later frequency difference (0.32 GHz) is much larger than the former (0.15 GHz). A similar tendency is also found at the experimental data for the AW array with \( \Delta w = 40 \) nm and the homogenous width arrays with \( w = 200 \) and 240 nm.

Interestingly, the frequency difference between AW arrays with \( \Delta w = 40 \) nm and the homogenous width array with \( w = 240 \) nm is 0.25 GHz. It is smaller than the frequency difference between the data for the AW array with \( \Delta w = -40 \) nm and for the homogenous width array with \( w = 200 \) nm (0.32 GHz).

### B. Antiferromagnetic Ground State

The AFM ground state [see MFM image in Fig. 4(e) and its explanation in Fig. 4(d) II] exists in a certain range of applied fields [Fig. 4(b)] within the major hysteresis loops for the arrays. It can be also stabilized in a range of fields on respective minor hysteresis loops.

Shown in Figs. 6(a)–6(b) are FMR absorption spectra measured while running minor hysteresis loops for AW arrays with \( \Delta w = 340 \) nm and \( \Delta w = 40 \) nm, respectively. These spectra were recorded while decreasing \( H_{\text{app}} \) from \( H_{\text{max}} \) to \( -H_{\text{sat}} \). A discontinuous change in frequency is observed for \( H_{\text{app}} = -100 \) Oe and \(-227 \) Oe, respectively, which corresponds to returning to the original FM order through switching of the wider NWs back to the \(-Y\) direction. One mode dominates the spectra in Figs. 6(a) and 6(b) in the range of existence of the AFM state. Interestingly, for \( \Delta w = 340 \) nm, the mode frequency is a monotonic function of the applied field, while for \( \Delta w = 40 \) nm, the mode has a nonmonotonic character with a maximum frequency at \( H_{\text{app}} = 50 \) Oe.

Figure 6(c) summarizes the field dependencies of the FMR frequency in the region of the AFM order for different \( \Delta w \) values. In this figure, we keep one wire width constant \( (w_0 = 200 \) nm) and vary the width of the second wire on the structure period \( w_0 \) from 160 nm to 540 nm. We find that for \( \Delta w > 0 \) the FMR frequencies have a tendency to converge to one point

![FIG. 5.](image-url)
at the largest fields for the AFM order. Interestingly, the field dependence for $\Delta w = -40$ nm has a different asymptote for the larger applied fields. As we will show below, this is connected to the fact that for $\Delta w < 0$ and $\Delta w > 0$ the widths of the narrower wire are different (200 nm for $\Delta w > 0$ and 160 nm for $\Delta w < 0$). This suggests that the acoustic mode for $\Delta w < 0$ and $\Delta w > 0$ dominates across the sample volume except for the areas near element edges. Furthermore, for symmetry reasons the mean values across the sample volume of $M_x$, $M_y$, $H_x$, and $H_y$ vanish. For this reason in the following analysis, we account for only the longitudinal components $M_x$ and $H_x$ and, consequently, for the dynamic components perpendicular to them ($m_y$, $m_z$, $h_y$, and $h_z$). We also, as usual, linearize the magnetic torque equation

$$\frac{i \omega m(x,y,z)}{\gamma} = -[\{M(x) + m(x,y,z)\} [H(x) + h(y,z)]],$$

(4)

where $\omega$ is the resonance frequency, $M$ and $H$ are the vectors of static magnetization and of static internal field, and $m$ and $h$ are the dynamic magnetization and the dynamic effective field in the structure. For our case of 3D confinement the internal static field $H$, $M$ and, consequently, $m$ and $h$ have all three spatial components. However, the static components along the applied field (the $x$-component) dominate across the sample volume except for the areas near element edges. Furthermore, for symmetry reasons the mean values across the sample volume of $M_x$, $M_y$, $H_x$, and $H_y$ vanish. For this reason in the following analysis, we account for only the longitudinal components $M_x$ and $H_x$ and, consequently, for the dynamic components perpendicular to them ($m_y$, $m_z$, $h_y$, and $h_z$).

The minor loop FMR data for the homogenous NW array with $w_0 = 200$ nm is shown in Fig. 6(d). We found that it is difficult to achieve a good AFM order in a homogenous NW array. The cross-shape mode profile indicates that the NW array is switching by forming FM order clusters with different magnetization direction. Two MFM images are shown as an inset to this figure. They confirm the cluster mechanism of magnetization switching for this NW array.

Shown in Fig. 6(e) are FMR traces for an array with $\Delta w = 40$ nm taken at remanence. One sees that the resonance frequencies at remanence for the FM ($f_{0,FM}$) and AFM ($f_{0,AFM}$) magnetic orders are different. A frequency difference $\Delta f_0$ is about 200 MHz, which is easily measurable with the broadband FMR.

Figure 6(f) displays the dependencies of $f_{0,FM}$, $f_{0,AFM}$, and $\Delta f_0$ on $\Delta w$. We observed that the resonance frequency for the FM state is always larger than for the AFM one. Furthermore, a clear trend is seen that $\Delta f_0$ decreases as the $\Delta w$ is increased.

The macroshape of the wire array is a rectangle $10 \mu m \times 20 \mu m$ [Fig. 1(a)] with the static magnetization vector directed along the shorter rectangle size. This produces a noticeable collective static demagnetizing field along the wires (i.e., along the $x$-direction). This also modifies the collective dynamic demagnetizing field with respect to previously studied case4 of wires of practically infinite length with respect to the structure period. For this reason the previously developed two-dimensional (2D) numerical simulation approach does not provide accurate quantitative description of the collective magnetic dynamics in our case; however, qualitatively predictions by this previous theory still should be valid. A three-dimensional (3D) numerical model is required to accurately simulate the experimental results. Construction of such a model and generally much more time-consuming simulations of magnetic dynamics in 3D are beyond the scope of the present paper. Instead in the following we suggest a simpler phenomenological approach to describe the dynamics for the AFM ground state. The approach is based on the idea of resonances of two sublattices of the artificial crystal, which are coupled by the collective dipole field.

The starting point of this approach is the linearized Landau-Lifshitz magnetic torque equation of motion for magnetic moment, which we cast in the following form:

$$\frac{i \omega m(x,y,z)}{\gamma} = -[\{M(x) + m(x,y,z)\} [H(x) + h(y,z)]],$$

(4)

where $\omega$ is the resonance frequency, $M$ and $H$ are the vectors of static magnetization and of static internal field, and $m$ and $h$ are the dynamic magnetization and the dynamic effective field in the structure. For our case of 3D confinement the internal static field $H$, $M$ and, consequently, $m$ and $h$ have all three spatial components. However, the static components along the applied field (the $x$-component) dominate across the sample volume except for the areas near element edges. Furthermore, for symmetry reasons the mean values across the sample volume of $M_x$, $M_y$, $H_x$, and $H_y$ vanish. For this reason in the following analysis, we account for only the longitudinal components $M_x$ and $H_x$ and, consequently, for the dynamic components perpendicular to them ($m_y$, $m_z$, $h_y$, and $h_z$). We also, as usual, linearize the magnetic torque equation

$$\frac{i \omega m(x,y,z)}{\gamma} = -[\{M(x) + m(x,y,z)\} [H(x) + h(y,z)]],$$

(4)

where $\omega$ is the resonance frequency, $M$ and $H$ are the vectors of static magnetization and of static internal field, and $m$ and $h$ are the dynamic magnetization and the dynamic effective field in the structure. For our case of 3D confinement the internal static field $H$, $M$ and, consequently, $m$ and $h$ have all three spatial components. However, the static components along the applied field (the $x$-component) dominate across the sample volume except for the areas near element edges. Furthermore, for symmetry reasons the mean values across the sample volume of $M_x$, $M_y$, $H_x$, and $H_y$ vanish. For this reason in the following analysis, we account for only the longitudinal components $M_x$ and $H_x$ and, consequently, for the dynamic components perpendicular to them ($m_y$, $m_z$, $h_y$, and $h_z$). We also, as usual, linearize the magnetic torque equation

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where $\omega$ is the resonance frequency, $M$ and $H$ are the vectors of static magnetization and of static internal field, and $m$ and $h$ are the dynamic magnetization and the dynamic effective field in the structure. For our case of 3D confinement the internal static field $H$, $M$ and, consequently, $m$ and $h$ have all three spatial components. However, the static components along the applied field (the $x$-component) dominate across the sample volume except for the areas near element edges. Furthermore, for symmetry reasons the mean values across the sample volume of $M_x$, $M_y$, $H_x$, and $H_y$ vanish. For this reason in the following analysis, we account for only the longitudinal components $M_x$ and $H_x$ and, consequently, for the dynamic components perpendicular to them ($m_y$, $m_z$, $h_y$, and $h_z$).
small-signal model. The total dynamic demagnetizing field $\mathbf{h}$ is the sum of dynamic self-demagnetizing field of the sublattice $W_1$ and the collective dynamic stray or coupling field from the sublattice $W_2$. Thus $h_z = -4\pi N_{xy1}m_1 - 4\pi N_{yx1}m_2$. Similarly, $h_x = -4\pi N_{yx1}m_1 - 4\pi N_{xy1}m_2$. Note that we do not include the antidiagonal components of the collective stray field tensor $N_{yx1}$, $N_{xy1}$, etc. Neglecting them is justified by the result in Ref. 6: the spatial mean values of the contributions to the total demagnetizing field they produce vanish for symmetry reasons.

The same analysis applies for the dynamic dipole field inside the wires of the sublattice $W_2$. The relevant tensors in this case are $N_{22}$ and $N_{21}$. Furthermore, we can use the same approach to describe static demagnetization. The static internal field $H$ can be written as $H = H_{app} + H_1$. One important component of $H_1$ is the self-demagnetizing field ($-4\pi N_{xx1}M_x$). There will be also contribution from the other sublattice. It useful to total the demagnetizing field using two different ways: $H_1 = -4\pi N_{xx1}M_x - 4\pi N_{xx2}M_x$ and $H_2 = -4\pi N_{xx2}M_x f_1$, where $f_1$ is a correction factor. (In these expressions the signs of $M_x$ and $M_z$ are given by the directions of static magnetizations of sublattices with respect to the applied field.) Thus $H = H_{app} + H_1$.

It is easy to infer that $f_1$ should be bigger than 1 for the FM ground state and smaller than 1 for the AFM. Again, the same analysis applies for $W_2$: $H = H_{app} + H_2$. $H_2 = -4\pi N_{xx2}M_x f_2 - 4\pi N_{xx2}M_x f_1$ and $H_2 = -4\pi N_{xx2}M_x f_2$. We also take into account Eq. (2) for the self-demagnetizing factors. As a result Eq. (4) transforms into

$$\left(\begin{array}{ccc}
\frac{-i\omega}{\gamma} & -(A_1 + v_{12}) & 0 \\
B_1 + v_{12} & \frac{-i\omega}{\gamma} & B_{12} \\
0 & -A_{21} & \frac{-i\omega}{\gamma} \\
B_{21} & 0 & B_2 + v_{21} & \frac{-i\omega}{\gamma}
\end{array}\right) \times \left(\begin{array}{c}
m_{y1} \\
m_{z1} \\
m_{y2} \\
m_{z2}
\end{array}\right) = 0$$

(5)

with $A_{1(2)} = H_{app} - 4\pi \Delta N_{xx1(2)}M_x^{1(2)}$, $B_{1(2)} = H_{app} - 4\pi \Delta N_{xx1(2)}M_x^{1(2)}$, $v_{12(21)} = 4\pi N_{xx2(1)}M_x^{1(2)}$, $A_{1(2)} = 4\pi N_{xx1(2)}M_x^{1(2)}$, and $B_{1(2)} = 4\pi N_{xx1(2)}M_x^{1(2)}$. For the FM order $M_{1x} = M_x$, and for the perfect AFM order $M_{1x} = -M_x$, $M_{1z} = -M_z$.

It is possible to make further important simplifications. First, because in the following we will apply this theory for the array with a small $\Delta w$ (40 nm) just to decrease the number of fitting parameters in the future fits, we assume that $N_{y1z} = N_{y1z}/N_{zyc}$ and $N_{x1z} = N_{x1z}/N_{zxc}$.

The nontrivial solution for Eq. (5) is given by the condition of vanishing of the determinant $|D|$ of the matrix of coefficients of Eq. (5). This results in a new equation $|D(\omega_0)| = 0$ for the unknown variable $\omega_0$. This new equation can be cast in a very simple form which is standard in the theory of coupled oscillators:

$$\left(\omega_1^2 - \omega_0^2\right)\left(\omega_2^2 - \omega_0^2\right) = \omega^2 C + E$$

(6)

with $\omega_1^2 = \gamma^2(A_1 + v_{12(21)})B_{12(21)}$, and $\omega_2^2 = \gamma^2(A_1 B_{21(22)} + A_{21(22)})$, $C = \gamma^2(A_1 B_{21(22)} + A_{21(22)})$, and $E = \gamma^2(A_1 + v_{12(21)})B_{12(21)}$. As one sees from Eq. (5), the quantities $\omega_1$ and $\omega_2$ may be obtained by equating the determinants of the $2 \times 2$ matrices, which represent the upper left-hand and lower right-hand $2 \times 2$ blocks of the $4 \times 4$ matrix in Eq. (5). In the absence of dipole coupling between the two sublattices ($N_{zyc} = N_{zxc}$), the right-hand side of Eq. (5) vanishes. Furthermore, for the uncoupled case, $v_{12(21)} = 0$. One then sees that $\omega_1$ and $\omega_2$ represent eigenfrequencies of the uncoupled sublattices.

As usual in the theory of coupled oscillators, from (6), one easily sees that when the frequencies of the uncoupled sublattices differ noticeably ($|(C + E/\omega_1^2)/\omega_1^2| \ll 1$), the eigenfrequencies $\omega$ of the two coupled states should coincide with the eigenfrequencies of the uncoupled sublattices. Only when the applied field is such that the frequencies of the sublattices are degenerate $\omega_1 = \omega_2$ or close to the degeneracy (i.e., when $|(C + E/\omega_1^2)/\omega_1^2| < 1$), the coupling removes the degeneracy to produce two frequencies noticeably different from $\omega_1$ and $\omega_2$. One of these solutions represents the acoustic mode $\omega_1$ and the second the optical mode $\omega_2$ of the coupled oscillations.

We employ the least-squares-fit method to fit our experimental data for the FM and AFM ground states with Eq. (6). We perform this fitting for the array with $\Delta w = 40$ nm and $s = 80$ nm. This fitting allows us to extract values of coupling demagnetizing factors from the experimental data. For these fittings we need data for uncoupled sublattices: $\omega_1$ and $\Delta N_{zyc,0}$, $\alpha = 1, 2$. They have been actually experimentally obtained in Sec. III, where we consider arrays with $s = 1.5 \times w$. Indeed, the distance between $W_2 (w = 200$ nm) on this AW array is $200 + 80 = 360$ nm, which differs just by 20% from $1.5 \times w$. Similarly, the distance between the wires $W_2$ on the AW array is $400$ nm, which differs by just 10% from $360$ nm = $1.5 \times w$.

The fitted dependencies for both FM and AFM magnetic orders are shown in Fig. 7 along with the data for the respective sparse arrays from Sec. III. One clearly sees that the model of two coupled sublattices correctly describes the behavior of the experimental dependencies. Indeed the mean deviation of the fitted curve from the experimental data is 0.055%.

The extracted values of demagnetizing factors are as follows. For the FM order: $N_{zyc} = 0.0118$, $N_{zxc} = 0.0921$, $N_{x1z} = 0.00165$, and $N_{x2z} = 0.000904$. For the AFM ground state: $N_{zyc} = 0.0032$, $N_{zxc} = -0.1915$, $N_{x1z} = -0.00165$, and $N_{x2z} = -0.00904$. These values are less than 25% of the self-demagnetizing factors ($\Delta N_{zyc,1} = 0.06563$, $\Delta N_{zyc,1} = 0.81206$, $\Delta N_{zyc,2} = 0.07331$, and $\Delta N_{zyc,2} = 0.78464$).

The important point in Fig. 7 is that the experimental and theoretical dependencies for the acoustic (lower-frequency) AFM mode and for the sparse array of the wide wires converge for large negative fields. One also sees that the theoretical optical (upper-frequency) AFM mode approaches the experimental data for the sparse array of the wide wires for large positive fields. This is in full agreement with our analysis above of Eq. (6) as an equation describing the dynamics of coupled oscillators.
Interestingly, there is no such striking convergence of the experimental data for the isolated narrow wires with the acoustic AFM mode for large positive fields and with the optical AFM mode for large negative fields, although these dependences clearly come close together. Even in the absence of the latter convergence, one can conclude that the nonmonotonic dependence of the FMR frequencies for the AFM states and formation of the forbidden frequency zone between the optical and the acoustical modes is due to hybridization of the resonances in the isolated wires. The hybridization results in repulsion of the falling and growing lines for the uncoupled wires in Fig. 7, which forms the prohibited zone.

The mode repulsion as the origin of the nonmonotonic field dependence of the resonance frequencies for the AFM state was first suggested in Ref. 35 based on an analysis of the simulation data obtained with the rigorous 2D numerical model for wires of infinite length. Our fit here using Kittel's equation. The red squares are minor-loop experimental results for an array of dipole-coupled AW NW array consisting of wires of the same width (w_1 = 240 nm and w_2 = 200 nm). The red lines are the calculated dispersion for the FM ground state (thin solid line) and for the AFM one (thick solid line is for the acoustic mode, thick dash line is for the optical mode). The thick and thin red dash-dotted lines are calculations for v_{12} = v_{21} = 0.

FIG. 7. (Color online) The black triangles and the blue circles are the experimental FMR dispersion of isolated NW array with w = 240 nm and w = 200 nm, respectively. The thin extended lines are the fittings using Kittel's equation. The red squares are minor-loop experimental results for an array of dipole-coupled AW NW array consisting of wires of the same width (w_1 = 240 nm and w_2 = 200 nm). The red lines are the calculated dispersion for the FM ground state (thin solid line) and for the AFM one (thick solid line is for the acoustic mode, thick dash line is for the optical mode). The thick and thin red dash-dotted lines are calculations for v_{12} = v_{21} = 0.

The physical reason for the difference between the values for N_{xx,12} and N_{xx,21} is just the wire width. One may expect that the largest contribution to the total collective static stray field originates from these narrow wires. Similarly, the largest contribution to the total collective static stray field at the positions of the wide wires originates from the two closest narrow wires. Each wide wire generates a larger stray field than each narrow wire. Furthermore, the distance between the wide wires on the NW array w_2 + 2s is smaller than the distance between the narrow wires w_1 + 2s. This suggests that the joint effect of these two contributions should result in a noticeably stronger collective static stray field at the positions of narrow wires than at the positions of the wide wires. This explains the striking difference in asymptotes.

The same reasoning applies for the FM state: the smaller fundamental mode frequency for the FM state in Fig. 7 than for those of the uncoupled stripes of both widths originates from both an increase in the internal static field and modifications in the dynamic field due to dipole coupling. Interestingly, the extracted values for N_{yy,12} and N_{yy,21} show that the coupling via the out-of-plane component of the dynamic field plays a noticeably larger role for the AFM ground state than for the FM state. This probably reflects the fact that for the AFM state magnetization precession is in difference senses in the wide and narrow wires.

For the latter reason for the acoustic AFM mode, both in-plane and out-of-plane components of the total internal dynamic field of wires decrease due to dipole coupling. For the same sense of precession in both types of wires, as is the case for the FM state, the out-of-plane dynamic demagnetizing field is increased due to coupling, whereas the in-plane one is decreased. Thus, this difference in precession senses explains the larger frequencies for the FM mode than for the AFM mode seen in Fig. 6(e). For H_{app} = 0, the difference in the frequencies for the FM and AFM magnetic orders originates completely from the difference in the demagnetizing fields; therefore, the data in Fig. 6(f) for the remanent state are the most illustrative. (Note that it follows from Eq. (5) that the collective static stray field partially masks this behavior: it increases the frequency of the AFM mode and decreases the frequency of the FM mode. As a result, the net difference between the
resonance frequencies for the FM and AFM states is smaller than expected from the purely dynamic coupling effect.)

Let us now come back to Fig. 5(a). The fundamental resonance frequency of uncoupled oscillations grows with decrease in the wire width [Fig. 2(b)]. As for any coupled system consisting of oscillators with two different eigenfrequencies in the absence of coupling, the acoustic and the optic modes of the coupled state are formed based on the respective eigenoscillation of one of these uncoupled oscillators. For instance, if the dynamic coupling reduces the eigenfrequency of both oscillators (which is the case of the FM state), the lower-frequency collective mode will be formed on the basis of the uncoupled resonance with the lower eigenfrequency [in our case on the basis of the resonance in the sublattice W₁ (the wide wires)]. Similarly, the higher frequency collective mode will be formed on the basis of the higher-frequency uncoupled resonance (one in the sublattice W₂). This simple conclusion follows from the fact that there should be a smooth transition without degeneracy in frequencies to the limiting case of the vanishing coupling. Since for periodic wire arrays, the collective dynamic demagnetization reduces the eigenfrequency of the in-phase magnetization precession and increases it for the antiphase precession, the acoustic (in-phase) collective mode is formed on the basis of the uncoupled oscillation in the wide wires and the optic (antiphase) mode on the basis of that in the narrow wires. Therefore one can expect a difference between the behaviors of the eigenfrequency of the acoustic mode as a function of Δw for positive and negative Δw values. Indeed, a large difference in the distances in frequency from the curve for Δw = 0 to the curves Δw = -40 nm and Δw = 40 nm is seen in Fig. 5(a).

However, the argument above suggests that qualitatively the behavior should be opposite to one observed in this figure: one should expect the resonance frequency for Δw = 40 nm to be close to the resonance frequency for the homogenous width array with w = 240 nm and the resonance frequency for Δw = -40 nm to the one for the homogenous array with w = 200 nm. This idea is fully confirmed by a numerical simulation using the 2D model for the stripes of infinite length. Because from Fig. 5(b) one clearly sees that this is not the case in our experiment (the experimental data for Δw = 40 nm are closer to w = 200 nm rather than to w = 240 nm), we may infer that the difference between the expected behavior and the measured one is due to noticeable variation in static demagnetization with Δw for our array of finite length.

VI. CONCLUSION

We have carried out a systematic experimental investigation of the dynamic responses of the NW arrays with different wire configurations using broadband FM resonance spectroscopy. Homogenous and AW Ni₈₀Fe₂₀ NW arrays of various geometrical parameters have been characterized.

The FMR responses of homogenous NWs are markedly sensitive to wire widths and interwire spacing due to a spatially varying demagnetizing field. The collective spin-wave mode profile for FM and AFM ordered ground states were investigated by controlling the relative alignment of magnetization of neighboring NWs. We found that the magnetic ground states of coupled AW NW arrays can be tuned by manipulating the magnetic field histories. We have imaged directly the FM and AFM ground state magnetic configurations using MFM. Simple analytical models, easy to use by experimentalists in the field, have been suggested to explain the experimental data. The theory evidences the nature of the collective resonance modes of the AFM ground state as originating from hybridization of eigenmodes of two antiferromagnetically coupled sublattices and from splitting of the respective branches of the dependencies of resonant frequency vs resonant field caused by this hybridization.

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