A nanomagnetic polarization amplifier for graphene

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The generation of nonequilibrium electron spin polarization, spin transport, and spin detection are fundamental in many quantum operations. We demonstrate that a lattice of magnetic nanodots enhances the electron spin polarization in monolayer graphene. We probed the polarization through a resistively-detected variant of electron spin resonance (ESR) and observed that the resonance amplitudes are amplified by the presence of the nanodots. Each nanodot locally injects a surplus of spin-polarized carriers into the graphene, and the ensemble of all these ‘spin hot spots’ generates a nonequilibrium polarization in the graphene layer at macroscopic lengths scales whenever the interdot distance is comparable or smaller than the spin diffusion length.

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The lifetime of a spin information in a two-dimensional carrier system can vary considerably between a few picoseconds and several microseconds, depending on the host material, imposed confinements, the temperature and various types of interactions\(^1\). Carrier transport with long spin lifetimes is key to store, transport, and compute quantum information using the electron spin. The low intrinsic spin-orbit interaction in monolayer graphene, the low atomic number of its carbon atoms and a very small hyperfine coupling between the nuclear and electron spins theoretically permits long spin lifetimes in the range of microseconds\(^2\)–\(^7\). In real graphene devices, however, the lifetime is limited by the interaction with the substrate and/or adatoms, which contribute extrinsic spin-orbit coupling and enhance parasitic spin relaxation mechanisms\(^8\). Experimentally determined spin lifetimes are thus found to be two orders of magnitudes smaller than theoretically anticipated\(^9\)–\(^21\). In devices fabricated from graphene synthesized by chemical vapor deposition (CVD) on a SiO\(_2\) substrate, we are confronted by a multitude of intrinsic and extrinsic defects that promote spin flips and govern spin relaxation.

In this work, we address the question whether CVD graphene can be refined to maintain a high (non-equilibrium) spin polarization over macroscopic distances by "doping" the supporting substrate with a lithographically defined hexagonal lattice of magnetic nanodots, with an average lattice constant comparable to the spin diffusion length, \(\lambda_s\). In electron spin resonance (ESR) experiments under a magnetic field perpendicular to the sample, and at low temperatures, we observe that graphene in this nanomagnet device experiences an enhancement of the electron spin polarization. We verify the enhanced polarization by analyzing the resonance peaks and comparing them to a reference sample without nanodots. The magnetic nanodots appear to act as spin hot spots that regenerate the spin polarization in the nearby 2D carrier system. Nanomagnetic/graphene hybrid structures may thus constitute an alternative strategy for controlling the spin polarization in materials with strong spin diffusion.

We created the (quasi-hexagonal) nanodot lattice in a multistep process by coating a self-assembled monolayer of SiO\(_2\) nanoparticles on a 3 nm Pt/0.7 nm Co/7 nm Pt stack on top of a SiO\(_2\)/p-Si substrate. The SiO\(_2\) particles act as shadow masks during ion milling, which leaves behind an array of cylindrical Pt/Co nanodots\(^22\). Monolayer graphene is transferred onto the substrate and structured into a large Hall bar of width \(W = 22 \mu\text{m}\) and 200 \(\mu\text{m}\) length (lateral voltage probe separation \(L = 100 \mu\text{m}\)) as detailed in Ref. 23. This is sample B. A reference sample (sample A) was prepared using the same source of CVD graphene and the same clean room processes; for this sample, the substrate was a pure SiO\(_2\)/p-Si wafer.
FIG. 1. Reference sample A (graphene on plain SiO$_2$). Upper panel (a): $\rho_{xx}^{\text{dark}}$ measured without microwaves (blue solid line) and $\rho_{xx}^{23 \text{ GHz}}$ under continuous radiation of 23 GHz/$+21 \text{ dBm}$ (red solid line). The back gate was tuned to the smallest density at the charge neutrality point of approximately $n \approx 2 \times 10^{11} \text{ cm}^{-2}$. Thermal activation leads to an overall increase of the conductivity under microwave radiation. Lower panel (a): $\Delta \rho_{xx} = \rho_{xx}^{23 \text{ GHz}} - \rho_{xx}^{\text{dark}}$. The peak centered at $B = 0$ T is resulting from thermal activation in the weak localization regime; the highlighted areas indicate electron spin resonances (see main text). (b) The resonance amplitudes in ohms, taken from Lorentzian fits of the outer (high field) resonances that occur around $0.85 \text{ T}$ as function of the carrier concentration. The black solid line represents $\rho_{xx}$ measured at $B = 0$ T by sweeping the gate voltage (right-hand axis). A discontinuity exists because the density can not be tuned to be zero for $T > 0$ K. The dashed red line represent the estimated amplitudes based on the Fermi distribution function (see main text for details). (c) Spin diffusion length ($\lambda_S$) versus frequency measured at $2 \times 10^{11} \text{ cm}^{-2}$.

Both samples were cooled down to nominally 1.3 K in the same custom-made vacuum probe station that is submerged in a liquid helium variable temperature insert (VTI). The cryostat is equipped with a superconducting magnet that generates a magnetic field ($B$) perpendicular to the sample plane. We employ a standard lock-in method that passes a low frequency alternating current of $I = 2 \text{ nA}$ and 37 Hz through the Hall bar. The lock-ins detect the longitudinal voltage drop, $V_{xx}$, and the Hall voltage, $V_{xy}$. The resulting longitudinal and Hall resistivities are calculated as $\rho_{xx} = \frac{W}{L} \cdot \frac{V_{xx}}{I}$ and $\rho_{xy} = \frac{V_{xy}}{I}$, respectively. At low temperatures and at the charge neutrality point, sample A has an (electron) density of $n \approx 2 \times 10^{11} \text{ cm}^{-2}$, a mobility of $\mu \approx 3200 \text{ cm}^2 \cdot (V \cdot s)^{-1}$ and a carrier mean free path for ballistic transport of $l_e = \frac{\hbar}{2\mu} \sqrt{\frac{T}{\pi}} \approx 16.7 \text{ nm}$. Sample B is characterized by an intrinsic hole density of $p \approx 1.5 \times 10^{13} \text{ cm}^{-2}$, a mobility of $\mu \approx 70 \text{ cm}^2 \cdot (V \cdot s)^{-1}$ and $l_e \approx 3.2 \text{ nm}$. The samples were exposed to microwaves through a two-turn coil.
located next to the sample. The coil is connected through a semi-rigid coaxial wire to a frequency generator. Microwaves are applied as continuous wave (CW) with constant frequency and constant power. The sample temperature under microwave radiation may increase up to 20-30 K.

The realm of the electron spin and the spin polarization in the two samples is experimentally accessible through a resistively-detected variant of electron spin resonance (RD-ESR). In a magnetic field the degeneracy between spin-up and spin-down electrons in the graphene is lifted and a Zeeman gap opens that is proportional to $B$. We probe this process and related spin properties by exposing a sample to microwaves of frequency $\nu$. Whenever $h\nu$ (h: Planck constant) matches the energy of the Zeeman gap $g \cdot \mu_B \cdot B$ ($g$: g-factor, $\mu_B$: Bohr magneton), the resonant absorption that is accompanied by spin flips is detectable as peaks in the sample resistance.

In pristine/plain graphene samples on an insulating substrate, the resonance peaks are usually buried in the resistive background, requiring the subtraction of a measurement without microwaves $\rho_{xx}^{\text{dark}}$ and the calculation of $\Delta \rho_{xx} = \rho_{xx}^{\nu} - \rho_{xx}^{\text{dark}}$. Figure 1 (a) illustrates this analysis using an exemplary data set measured on our reference sample A for $n \approx 2 \times 10^{11}$ cm$^{-2}$. The calculation of $\Delta \rho_{xx}$ reveals two resonance peaks axially symmetric with respect to $B = 0$. The prominent outer high field peaks highlighted in red represent the resistive response to resonant spin-flips when $h\nu$ matches the (pure) Zeeman energy. The weak inner ESR peaks (grayed) are related to the intrinsic spin-orbit coupling band gap in graphene. In the following, we focus our analysis on the prominent outer peak (see Supplementary Material).

We will now scrutinize the resonance line shapes and amplitudes, which encode information on spin diffusion and spin polarization. The Lorentzian line width of the resonance peaks ($\Delta B_{\text{res}}$) is proportional to the spin diffusion length, $\lambda_S \propto \Delta B_{\text{res}}^{-0.5}$ (see Supplementary Material). Figure 1(c) shows $\lambda_S$ measured on sample A deduced at various frequencies, resulting in a mean value of $\lambda_S = (346 \pm 60)$ nm. The amplitude of the microwave-induced resistance peak, on the other hand, is a measure of the transition probability between two spin states. For resistive detection of spin resonances, a sufficient number of electrons must be resonantly excited from an equilibrium spin ground state and flip their spins. This probability is dictated by the number of available initial and final states, or more precisely, by the spin polarization that maps the difference between the number of spin-down $N_\downarrow$ and spin-up $N_\uparrow$ electrons.

Figure 1 (b) shows the evolution of the peak height with carrier concentration and type for sample A (for a constant magnetic field, i.e., for a constant resonance frequency of 23 GHz). The resonance amplitude is largest around the charge neutrality point (CNP), where it culminates at 150
Ohms for this particular microwave frequency and power, but rapidly decays at higher densities and eventually vanishes in the resistive background. At a hole concentration of $p \approx 1 \times 10^{12}$ cm$^{-2}$, the peak height has dropped to approximately 10 Ω. The monotonic asymptotic decay of the amplitude is the result of the increasing Landau level filling with density at low $B$, a regime in which only a fraction of electrons can be involved in the resonant absorption processes. We can estimate this behavior using the following model.

The thermal energy even under strong radiative heating never exceeds $k_B \cdot T = 2.5$ meV, which is still dwarfed by the Zeeman energy of about 100 meV. The resistive signal is therefore proportional to $f(E(\mu_F) - E_{\text{Zeeman}}/2)(1 - f(E(\mu_F) + E_{\text{Zeeman}}/2))$, where $f$ is the Fermi distribution function and $E(\mu_F)$ is the free energy. The free energy is of the order of 0.1 - 1 meV and depends linearly on the chemical potential $\mu_F$ (Supplementary Material of Ref. 26). The signal (or $f$) thus decays exponentially with density. We employed a mean field calculation to estimate the Fermi energies ($E_F = n \text{[cm}^{-2}\text{]} \cdot 0.771 \times 10^{-11}$ [meV]) and the expected resonance amplitudes that are proportional to $f(E_{\text{initial}}) \cdot (1 - f(E_{\text{final}}))$, with $(E_{\text{final}} - E_{\text{initial}})$ being equal to the Zeeman energy. The results are shown by the red dashed line in Fig. 1 (b), which reasonably well reproduces the overall trend of our data for the sample under radiative heating from microwaves radiation. At high carrier densities, the signal disappears as the majority of spin-unpolarized carriers mask all resonant absorption processes of unpaired/spin-polarized electrons.

By placing a sheet of monolayer CVD graphene over an artificial lattice of magnetic nanodots (sample B), we can introduce a surplus of spin-polarized carriers, which amplifies the resistive response under resonance. Figure 2(a) shows a scanning electron microscopic (SEM) image of our array of cylindrical Pt/Co nanodots on a SiO$_2$ substrate (sample B). The dots have an average diameter of 27 nm and a height of 10 nm. They form a low-order hexagonal lattice and cover 1.4% of the substrate surface with a mean separation of approximately 200 nm. We chose this separation because it is smaller than the typical spin diffusion lengths in our CVD graphene devices. Each nanodot represents a single ferromagnetic domain with its magnetization pointing out of plane$^{22,30}$ due to interface anisotropy. The magnetic stray field perpendicular to the nanodot surface reaches approximately 59 mT, however, at lateral distances of 200 nm the stray field is too small to (anti)ferromagnetically couple neighboring dots$^{31,32}$. The platinum in the capping layer of the nanodots possesses a very high electron affinity, which induces strong $p$-doping and the high intrinsic hole concentration in the graphene of sample B ($p \approx 1.5 \times 10^{13}$ cm$^{-2}$).

Figure 3 (a) shows two typical measurements of the now as measured (absolute) longitudinal...
FIG. 2. Sample B. (a) SEM image of the substrate with prepatterned Pt/Co nanodots. (b) Schematic illustration of the graphene Hall bar sample on the substrate and the relative orientations of the microwave field generated by a nearby coil, the external magnetic field ($B$), and the low frequency transport current ($I$).

resistivity on sample B under constant microwave radiation of frequencies $\nu = 27$ GHz and 35 GHz. Despite the very high intrinsic hole concentration, which is two orders of magnitude larger than for the data shown in Fig. 1 and still one order of magnitude larger than those densities for which the resonance peak had vanished in the resistive background, large resonances are seen. We stress that a similar density-dependent measurement as shown in Fig. 1(b) was not feasible since the nanodots appear to screen the electric field from the back gate; even excessively large voltages only marginally affected the carrier density. Though, based on the behavior seen in Fig. 1(b), we reason that the resonance amplitudes in sample B would be further enhanced at lower carrier concentrations. To estimate this hypothetical increase, we "reverse-engineered" the simulation for the resonance amplitudes using the results at $p_i$ that correspond to a certain surplus spin polarization. Close to the CNP at a density of $\approx 1-2 \times 10^{11}$ cm$^{-2}$, the predicted amplitude should be a factor of $23 \times$ larger than at $p_i$ [Fig. 4 (b)].

By comparing the resonance peak amplitudes of samples A and B, measured under the same experimental conditions (e.g., temperature, microwave power), we can stipulate a relative change in the ratio of $\frac{N_{\uparrow}}{N_{\downarrow}}$ that is induced by the magnetic nanodots. We stress that the resonance peaks
FIG. 3. Sample B. (a) As measured (i.e., absolute) $\rho_{xx}$ shown for two exemplary frequencies, $\nu = 27$ GHz and 35 GHz using a power of +21 dBm, shows strong Zeeman resonance peaks (red highlighted areas) and transitions originating from the existence of an intrinsic band gap (unmarked). Inset: Frequencies vs. resonance positions obtained from Lorentzian fits of the resonances. The slope corresponds to a $g$-factor of 1.96±0.02 and the intercept of the upper line to an energy of $(42.35\pm1.07)\mu$eV. All measurements were performed with the (intrinsic) hole density of $p_i \approx 1.5 \times 10^{13}$ cm$^{-2}$, which is two orders of magnitude larger as for the data shown in Fig. 1(a). (b) Ratio of peak heights of $\Delta_A$ (sample A at $n = 2 \times 10^{11}$ cm$^{-2}$) and $\Delta_B$ (sample B at $p_i$) vs. frequency.

exhibit a weak pseudo-dependence on the frequency that reflects the frequency-dependent attenuation of the microwaves in the coaxial wave guide, which can affect the coupling of microwaves to the electron spins and the sample temperature. The calculation of the ratios of the resonance peak height $\Delta_B$ measured in sample B (at its intrinsic density) and $\Delta_A$ measured in sample A (tuned to $n = 2 \times 10^{11}$ cm$^{-2}$) at matching frequencies eliminates these technological perturbations, as shown in Fig. 3(b). We find the mean ratio to be 15.2±4.4, i.e., the peaks in sample A are still larger by a factor of 15 due to the 100× smaller carrier concentration, however, the peak ratio is independent of other parameters. This implies that the nanodots have boosted the electron spin polarization in the graphene of sample B. We emphasize that not all measurements return a reasonable fit. While
we can always obtain the resonance position [inset Fig. 3(a)] from taking the derivative of the curve, we cannot always deduce a reliable amplitude; for that reason, not all frequencies measured in both samples can be compared.

To verify that we are not accidentally probing any parasitic effect in sample B, we evaluated the ESR frequency-dependence and plotted $\nu$ versus the occurrence of the resonance in $B$ [inset Fig. 3(a)]. The lower line represents the pure Zeeman gap, whereas the upper line is related to intrinsic spin-orbit (SOC) interaction that opens up a mini band gap in monolayer graphene$^{26}$ (see Supplementary Material). From a linear fit of the resulting dispersions, we can deduce the electron $g$-factor of $1.96\pm0.02$ and a SOC gap of $(42.35\pm1.07) \mu$eV. Both values confirm previous reports of ESR on graphene$^{24,26}$. Although each nanodot has a maximal stray field of approximately 59 mT, the total mean stray field of all dots is less than 1 mT averaged over the sample area. The stray field is thus too small to generate a detectable deviation from the previously reported $g$-factor of $1.95\pm0.01$, and neither does it play a role in the polarization enhancement (see Supplementary Material). We stress that the perseverance of the upper resonance line in the presence of magnetic perturbations underlines the existence of a topological phase of matter that is associated with the intrinsic SOC gap.

The observation of electron spin resonances with a $g$-factor of $1.96\pm0.02$ indicates that our measurements are probing the graphene layer and not the Pt/Co dots. The nanodots, however, manipulate the ESR transition probability in the graphene. We propose that the nanodots act as ‘spin hot spots’ that boost the polarization by carrier exchange, i.e., electrons enter the magnetic nanodots where they become spin-polarized before they are re-injected into the graphene. Carriers originating from the Fermi surface of cobalt have a spin polarization of about $20\%^{33,34}$. During the passage through the platinum, some of this polarization will be lost due to the small spin diffusion length in Pt$^{35}$. The polarized carriers emerging from such a spin hot spot will be subject to various sources for spin relaxation. However, as the spin diffusion length $\lambda_S = (220\pm49) \text{ nm}$ in sample B [inset Fig. 4(b)] is comparable to the distance between spin hot spots, an excess polarization is conserved before the electrons reach the next dot. These repeated processes can maintain higher (non-equilibrium) net spin polarization over macroscopic distances, which we observe as enhanced resonance peaks. Figure 4(a) schematically illustrates the underlying principle using two dots for two distinct regimes in which the spin diffusion is much larger and smaller than the interdot distance.

As a concluding remark, we emphasize that our system does not constitute a spin-valve device
or a spintronics application in the classical sense, having a specific source and drain spin configuration. Here, we outlined a possible route to amplify, attenuate and control the electron spin polarization over macroscopic distances in graphene by "doping" the supporting substrate with an array of magnetic Pt/Co nanodots that act as spin hot spots. We verified the enhanced spin polarization through resistively-detected electron spin resonance measurements. Utilizing a doping scheme for signal amplification has shown to be vital in optical signal transmission in fibers, for example, and might also proof to be useful in van-der-Waals materials.

FIG. 4. (a) Simplified schematic illustration of two contrasting diffusive regimes after carriers experience a polarization boost in the vicinity of a nanodot. For $\lambda_S/a \geq 1$ (top panel), a net excess spin polarization is maintained during the transit to the next 'spin hot spot', where the polarization is refreshed. For $\lambda_S/a \ll 1$ (bottom panel), the excess polarization is lost long before the carriers reach the next dot. Parameter $a$ represents the mean interdot distance. (b) The predicted enhancement factor of the resonance amplitude for lower carrier concentrations. Inset: exemplary data points for $\lambda_S$ versus $\nu$. The average spin diffusion length in sample B is smaller than in sample A due to the different $l_e$ that determines the spin diffusion constant and consequently $\lambda_S$. 

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SUPPLEMENTARY INFORMATION

DETERMINATION OF THE $g$-FACTOR

The electron spin resonance (ESR) line shape is broadened as a result of the exchange interaction between localized and conduction electrons and can be described by a Lorentzian function [Chemical Physics Letters 557, 118 (2013)]. From a Lorentzian fit of the resonance peaks we can extract the exact resonance positions as well as the half-width $\Delta B_{\text{res}}$ [the latter is needed for the determination of the spin diffusion length, see section below]. Figure S5 shows the determined resonance frequencies as a function of the magnetic field ($B$). The lower line represents the pure Zeeman gap, whereas the upper line is related to intrinsic spin-orbit (SOC) interaction that opens up a mini band gap in monolayer graphene [Physical Review Letters 122, 046403 (2019)].

We note that not all measurements return a reasonable fit. In these cases, we can still obtain the resonance positions (but not the half-widths) from taking the derivative of curve.

![Graph showing resonance frequencies as a function of magnetic field](image)

**FIG. 5.** ESR peak positions obtained from Lorentzian fits on sample A. ESR was measured at 1.4 Kelvin with the gate voltage tuned to 18 V, or $\Delta V_{\text{CNP}} = +1.5$ V, respectively. Linear fits yield an electron $g$-factor of $1.94 \pm 0.01$ from the slope as well as the intrinsic gap energy of $(42.78 \pm 0.99) \mu\text{eV}$ from the intercept of the upper line, which agrees with previous reports.
Our analysis in the main text is based on the (pure) Zeeman peak since these resonances are generally more pronounced and thus easier to fit. However, both resonances exhibit the same dependence on the carrier concentration as shown in Fig. S6.

FIG. 6. Sample A. The resonance amplitudes in ohms, taken from Lorentian fits of the pure Zeeman resonance and the resonances originating from the SOC gap as function of the gate voltage difference from the charge neutrality point (CNP). The data were obtained for 23 GHz and +21 dBm. The red solid line represents $\rho_{xx}$ measured at $B = 0$ T by sweeping the gate voltage (right-hand axis).
Note: We define the resonance amplitude (in the main text symbolized by $\Delta_{A/B}$) from a baseline in the resistive background from which we take the difference to the Lorentzian peak. An example is shown in Fig. S7. Errors are estimated based on the baseline and the accuracy of the fit. This, however, is not feasible for very small amplitude, or peaks overlaid by large background noise or slopes.

![Graph](image)

FIG. 7. Exemplary data shown as a black solid line. The red solid line is a Lorentzian fit.
DETERMINATION OF THE SPIN DIFFUSION LENGTH

We extract the spin diffusion length ($\lambda_S$) from the electron spin resonance experiments as follows: From the Lorentzian fit of a resonance peak, we obtain its half-width $\Delta B_{\text{res}}$, which is proportional to the spin relaxation time ($\tau_S$) [Nature Communication 3, 996 (2012)],

$$\tau_S = \frac{h}{4\pi \Delta E_{\text{res}}} = \frac{h}{4\pi \cdot g \cdot \mu_B \cdot \Delta B_{\text{res}}}.$$  

(1)

Here, $g$ is the g-factor of 1.94±0.01 as determined in section above. The corresponding spin diffusion length ($\lambda_S$) is given by

$$\lambda_S = \sqrt{D \tau_S}$$  

(2)

with the spin diffusion constant $D = 0.5 \cdot v_F \cdot l_e$. We use a Fermi velocity in graphene on SiO$_2$ of $v_F = 10^6$ m/s. [Physical Review B 78, 21408 (2008), Scientific Reports 2, 590 (2012)]. We use error propagation to include the errors from the $g$-factor and the width of the resonance peaks.

THE IMPACT OF STRAY FIELDS IN SAMPLE B

The main contribution for the enhanced degree of polarization in our device is originating from carrier exchange with the magnetic nanodots. To determine any possible impact from the dots’ stray fields, we make the following simplified estimate based on the Boltzmann distribution: the polarization depends on the external field $B$ and the stray field of approximately 60 mT directly over the dot. The positive $g$-factor of graphene dictates that electrons are in the energetically more favorable spin-down state from which they can get excited under microwave radiation. The Boltzmann distribution for the ratio of electrons with opposing spins is given by

$$\frac{N_\uparrow}{N_\downarrow} = \exp\left(\frac{-\Delta E_{\text{Zeeman}}}{k_B \cdot T}\right) = \exp\left(-\frac{g \cdot \mu_B \cdot B}{k_B \cdot T}\right).$$

(3)

with $N_{\uparrow,\downarrow}$ representing the number of spin up and spin down states, $k_B$ being the Boltzmann constant, $\mu_B$ the Bohr magneton and $\Delta E_{\text{Zeeman}}$ the Zeeman splitting, which is proportional to a magnetic field, $B$. The enhancement by the dots is dictated by their stray field, so that the increase in polarization by the nanodots can be expressed as a temperature-dependent factor that depends
only on the stray field:

\[
\frac{\binom{N_i}{N_i}}{\binom{N_i}{N_i} \text{ with dots}} = \exp\left( + \frac{g \cdot \mu_B \cdot B_{\text{stray field}}}{k_B \cdot T} \right)
\]  

(4)

We note that the graphene layer conforms to the nanodot lattice non-uniformly. As a consequence, the stray field that emanates from the dots penetrates the graphene under various angles and with various strength. The stray field strength itself may also vary and might be stronger in some dots. For that reason Fig. S8 provides an estimate for a range of stray fields.

FIG. 8. Estimate for a potential gain in polarization by the stray field based on the Boltzmann distribution.

Since only 1.4% of the sample area is covered by dots and since radiative heating increases the sample temperature, any gain induced by the stray fields should be negligible for the degree of spin polarization in the graphene layer under our experimental conditions. In ESR measurements, the dots’ stray fields will also modify the resonance frequency, however, also as a result of the low filling, this shift is below the detection threshold.
COMMENT ON STRAIN

The conformation of the graphene layer to the lattice of nanodots will certainly generate strain in the graphene. Since strain is linked to gauge fields, the natural question arises whether this could increase the spin polarization. We have studied electron spin resonance on monolayer graphene resting on modulated substrates with trenches of various heights and pitch, where strain-induced gauge fields would exceed those originating from the homogenously distributed nanodots. In all these samples, we never made a similar observations. Thus, we conclude that any strain in the graphene layer, arising from the presence of the nanodots, cannot account for the effects discussed in this manuscript.

COMMENT ON COMMENSURABILITY EFFECTS IN SAMPLE B

The longitudinal magneto resistance of a two-dimensional carrier system under the influence of a periodic array of artificial scatterers can show commensurability effects [Europhysics Letters 8, 179 (1989), Physical Review Letters 66, 2790 (1991)]. This phenomenon emerges when the magnetic field is strong enough to force electrons into performing (ballistic) orbits that match a certain geometry in a background potential. The nanodot array acts as such an artificial background potential that could localize electrons around a single dot or around an ensemble of dots. Commensurability effects have been observed as prominent peaks in the longitudinal resistance [Applied Physics Letters 93, 122102 (2008), Nano Letters 12, 8402 (2015)] in graphene. In our sample, however, the mean free path is only a few nm and thus too small for ballistic orbits.