Excitonic nature of magnons in a quantum Hall ferromagnet

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Magnons enable the transfer of a magnetic moment or spin over macroscopic distances. In quantum Hall ferromagnets, it has been predicted1 that spin and charge are entangled, meaning that any change in the spin texture modifies the charge distribution. As a direct consequence of this entanglement, magnons should carry an electric dipole moment. Here we report evidence of this electric dipole moment in a graphene quantum Hall ferromagnet1,5 using a Mach–Zehnder interferometer. As magnons propagate across the insulating bulk, their electric dipole moment modifies the Aharonov–Bohm flux through the interferometer, affecting both phase and visibility of the interference pattern. In particular, we relate the phase shift to the sign of this electric dipole moment and the loss of visibility to the flux of emitted magnons, and we show that the magnon emission is a Poissonian process. Finally, we probe the emission energy threshold of the magnons for transient states, between ν = 0 and ν = 1, and link them to the emergence of the gapless mode predicted in the canted-antiferromagnetic phase at charge neutrality6. The ability to couple the spin degree of freedom to an electrostatic potential is a property of quantum Hall ferromagnets that could be promising for spintronics.

The full spin polarization of electrons in a material is helpful for applications in spintronics, which led, for instance, to the development of a new type of ferromagnetic material, the so-called half-metallic alloys that show complete spin polarization in the 90s. It was demonstrated that a fully polarized system can also be obtained in two-dimensional materials in the quantum Hall effect regime. More precisely, in the N = 0 Landau level, the kinetic energy is frozen by the magnetic field and all the states are degenerate. At the Landau-level filling factor of ν = 1, the Coulomb energy E_C = e^2/(εl_B)^2 —much larger than the Zeeman energy E_z = g μ_B —plays a primary role and leads to a fully antisymmetric spatial part of the wave function. Therefore, the spin part is symmetric and the ground state is an ideal ferromagnet1, fully spin polarized and described by a simple Slater determinant. In the bulk, excitations of this quantum Hall ferromagnet are itinerant deformations of the uniform background magnetization known as magnons. These excitations are charge neutral and comprise an electron–hole pair or exciton separated by a distance l_ch = k μ_B (ref. 4), where k is the magnon wavevector that depends on the magnon energy and l_B = h/eB is the magnetic length. This feature relies on a unique property of a quantum Hall ferromagnet where the dynamics of spin and charge are entangled, and any modification of the spin texture leads to a change in the electronic density. Magnon generation is then necessarily tied to the propagation of an electric dipole moment.

Graphene, which shows a rich sequence of quantum Hall plateaus as spin and valley degeneracies are lifted in the low-lying Landau levels, is an ideal platform to study collective excitations at a filling factor of ν = 1 (refs. 10,11). Recent experiments12–14 have demonstrated magnon generation by an out-of-equilibrium occupation of edge channels. Magnon detection relies on absorption at the local vicinity of ohmic contacts. The threshold energy to excite these propagating collective modes is typically E_EB = g μ_B B. However, since detection relies on magnon dissipation in an ohmic contact, neither spin nor electrical dipole properties of these collective excitations have been addressed.

In this study, we reveal the spin–charge entanglement at ν = 1 in graphene by detecting the excitonic part of magnons. We first study the reflection properties of magnons at the interface of a p–n junction. We then tune such a p–n junction into a Mach–Zehnder interferometer (MZI)15–19 and measure the evolution of phase and visibility of interferences when the magnons are emitted. We show that magnons induce dephasing that cannot be accounted for by the magnetic moment carried by the magnon but only by its electric dipole moment. We also show that the loss of visibility can be accounted for by a very simple model, in which decoherence exponentially decreases with the Poissonian flux of the emitted magnons. Poissonian emission is an important observation since it supports the presence of magnons defined as elementary excitations. Other experiments10–14 cannot distinguish between an interpretation based on magnons or simple spin waves. We finally study the threshold energy emission of the magnons for transient states between the filling factors of ν = 1 and ν = 0, hinting at a transition towards the canted-antiferromagnetic phase of the ν = 0 state.

The sample is set in the quantum Hall effect regime by applying a strong perpendicular magnetic field; therefore, the current propagates along the edge channels, as shown in Fig. 1a. A top gate enables the tuning of the filling factor below it (denoted ν_T), the filling factor outside of the top-gated region is denoted ν_EB). Figure 1b plots the variation in the Hall resistance with the top-gate voltage and bias applied to the sample, allowing to precisely determine the filling factor in the sample and showing an absence of the breakdown effect at the biases used throughout this study. Due to contact doping, the local filling factor close to the ohmic contacts is ν ≥ 2 (ref. 13) (a dedicated discussion is provided in Supplementary Section VII. Applying a voltage V_E on the emitter contact (denoted E in Fig. 1a), a chemical potential difference μ = −eV_E is set between the inner and outer edge states (depicted in red and blue, respectively, in
When \(|eV_{\text{d}}| > E_{\text{g}}\) (\(E_{\text{g}} \approx 1\) meV at 9 T), electrons can tunnel from the inner to the outer edge channel, accompanied by a spin flip that generates a magnon\(^1\). The emitter sites directly depend on the sign of voltage \(V_{\text{E}}\); in Fig. 1a, these sites are denoted \(e\) for positive \(V_{\text{E}}\) and \(e^\text{−}\) for negative \(V_{\text{E}}\). The emitted magnons can propagate through the insulating bulk and be absorbed at a distant ohmic contact by the reverse process, where an electron tunnels from the outer to the inner edge channel accompanied by an inverse spin flip\(^1\). This leads to the emergence of a finite non-local signal \(\frac{dV_{\text{top}}}{dV_{\text{E}}}\), where \(V_{\text{top}}\) is the a.c. voltage applied (in addition to \(V_{\text{E}}\)) at the emitter contact and \(V_{\text{top}}\) is the non-local voltage on the \(i\)th detector contact.

Following ref. \(^1\), the absorption of a magnon on contact \(i=0, 1\) and 2 creates a chemical potential shift \(\varepsilon_{\text{A,i}}\) in the edge channel flowing from this contact. Similarly, the chemical potential shift generated by the absorption on the grounded contact upstream of contact \(i\) creates a chemical potential shift \(\varepsilon_{\text{G,i}}\) for \(i = 0\) or 2 (Fig. 1a). On contact 0, we thus have \(\frac{dV_{\text{top}}}{dV_{\text{E}}} = \left(\frac{d\varepsilon_{\text{A,0}}}{d\varepsilon_{\text{G,0}}} - \frac{d\varepsilon_{\text{A,2}}}{d\varepsilon_{\text{G,2}}}\right)\). We first study the non-local voltage properties as a function of \(\nu_{\text{E}}\) and \(\nu_{\text{top}}\). Figure 1c shows \(\frac{dV_{\text{top}}}{dV_{\text{E}}}\) as a function of the top-gate voltage \(V_{\text{top}}\) (changing \(\nu_{\text{E}}\) from \(-1\) to \(+1\)) at \(\nu_{\text{top}} = 1\). At \(\nu_{\text{top}} = 1\), the generation and propagation of magnons above \(V_{\text{top}} = 1\) meV is revealed by the finite non-local signal \(\frac{dV_{\text{top}}}{dV_{\text{E}}}\). This indicates that some of the emitted magnons propagate across the insulating bulk between the emitter contact and contact 0, and are absorbed on the latter. This observation is consistent with previous works\(^{12-14}\). Under the same conditions, we observe a clear non-local signal \(\frac{dV_{\text{top}}}{dV_{\text{E}}}\) on contact 1 (Fig. 1d), further demonstrating the propagation of magnons across the bulk. Note, however, that the non-local signal \(\frac{dV_{\text{top}}}{dV_{\text{E}}}\) on contact 2 is essentially zero; to understand this, we remind the reader that the non-local signal \(\frac{dV_{\text{top}}}{dV_{\text{E}}}\) also depends on the absorption of magnons on the nearest upstream grounded contact. In this case, the distance between contact 2 and the grounded contact is much smaller (~1 \(\mu\)m) than the distance from the emitter contact (~12.5 \(\mu\)m; Supplementary Fig. 5), so that they roughly receive the same amount of magnons, leading to a vanishingly small signal.

For \(\nu_{\text{top}} = 0\), \(\frac{dV_{\text{top}}}{dV_{\text{E}}}\) strongly increases (Fig. 1c), whereas \(\frac{dV_{\text{top}}}{dV_{\text{E}}}\) and \(\frac{dV_{\text{top}}}{dV_{\text{E}}}\) go to zero (Fig. 1d,e). This increase in \(\frac{dV_{\text{top}}}{dV_{\text{E}}}\) corresponds to the magnons reflected at the \((\nu_{\text{top}}, \nu_{\text{E}}) = (0, 1)\) interface, which are absorbed in contact 0 and cannot reach contacts 1 and 2. A possible explanation for this strong reflection would be that the mismatch of magnon velocities between \(\nu_{\text{top}} = 1\) and \(\nu_{\text{E}} = 0\) prevents magnons from being transmitted across the interface\(^1\). Interestingly, the non-local signal \(\frac{dV_{\text{top}}}{dV_{\text{E}}}\) remains constant as \(\nu_{\text{E}}\) is set from 0 to \(-1\), showing that magnons are still reflected on the \((\nu_{\text{top}}, \nu_{\text{E}}) = (-1, 1)\) interface. This observation is not universal\(^1\); we argue that it might stem from the formation of a thin strip at \(\nu = 0\) along the interface at \((\nu_{\text{top}}, \nu_{\text{E}}) = (-1, 1)\) that reflects magnons. The existence of this thin strip is clearly dependent on the sample, as it is affected by the sample electrostatics as well as disorder. Using simple conservation laws at the interface (details are provided in Supplementary Section VI), we can obtain an expression of the magnon transmission that is directly linked to their dispersion relation in each region. We thus obtain reasonable estimates of the spin stiffnesses at \(\nu = 0\) and \(\nu = 1\)—the central parameters of the quantum Hall ferromagnet model that quantify the ability to excite magnons.

At the transition between \(\nu_{\text{top}} = 1\) and \(\nu_{\text{E}} = 0\), \(\frac{dV_{\text{top}}}{dV_{\text{E}}}\) sharply increases (Fig. 1c), corresponding to a strong imbalance between the absorption on two consecutive contacts located close to one another. This indicates that the propagation of magnons is strongly non-ballistic in this regime; therefore, the magnons are deflected on length scales of the order of micrometres. One possibility would be that the deflection in the bulk is due to the excitonic nature of magnons. Indeed, magnons in the quantum Hall ferromagnet carry an electric dipole moment \(p = |e| \hbar \mathbf{z} \times \mathbf{k}\), where \(\mathbf{k}\) is the centre-of-mass momentum and \(\mathbf{z}\) is collinear to the magnetic field\(^6\) (Fig. 1a, yellow arrow). This electric dipole interacts with any charged localized states that...
typically form in the bulk of the sample, deflecting the magnons. However, at this point, we cannot completely rule out that this effect observed at the edge of the quantum Hall plateau may simply be due to small density inhomogeneities in the region near contact 2.

To explore the excitonic nature of magnons, we now tune the sample in a bipolar quantum Hall state to form an MZI, as described in our previous study. In the n-type region, the filling factor is $\nu_n = 1$ and one spin-polarized channel (Fig. 2a, blue arrow) circulates anticlockwise along the boundary, whereas in the p-type region, $\nu_p = -2$ and two channels with oppositely polarized spins circulate clockwise (Fig. 2a, red and yellow arrows). By using electrostatic side gates at the intersections between the physical edge of graphene and the p–n junction interface, we can realize two valley splitters that coherently split the two edge channels with the same spin (but opposite valley isospin) stemming from the n- and p-type regions (Fig. 2a, blue and yellow arrows), which then co-propagate along the p–n junction interface before recombining on the opposite side of the p–n junction.

We measure the interferometer’s electronic transmission probability $T_{MZ} = I_T/I_{inj}$ (where $I_{inj}$ is the injected a.c. current and $I_T$ is the transmitted current through the p–n junction measured at contact T; Fig. 2a) as magnons are generated. To do so, a d.c. bias $V_E$
is applied on the emitter contact E, which, importantly, is located downstream of the interferometer. Figure 2b shows the variations in $T_{\text{MZ}}$ as a function of the magnetic field and $V_e$.

At the threshold energy $|eV_e| \approx E_p$, we observe a clear phase shift in $T_{\text{MZ}}$ that increases with the applied d.c. bias, accompanied by a drop in the oscillation visibility $V$. This is explicitly illustrated in Fig. 2c, which shows that both phase shift and visibility sharply change at the threshold. The phase shift corresponds to a variation in the Aharonov–Bohm flux through the MZI; thus, the observed effect can only be due to the magnons changing the area of the MZI or the local magnetic field. One can rule out the latter effect by noting that for a slightly different value of the applied magnetic field, the phase shift is reversed. This is illustrated in Fig. 2d,e, which shows the same measurements as that in Fig. 2b,c, at an applied magnetic field of $B \approx 8.875$ T instead of $B \approx 8.850$ T. This change in sign of the phase shift cannot be attributed to the magnetic moment of the magnons, which is always of the same sign.

Instead, we argue that the decrease in both phase shift and visibility is due to the excitonic nature of magnons. Indeed, as depicted in Fig. 2a, depending on their incidence angle with respect to the $p$–$n$ junction, the dipole carried by the magnons acts as a local electrostatic field, affecting the area of the MZI and therefore its AB phase, as mentioned earlier. This phase is particularly sensitive to the projection of the magnon electric dipole perpendicular to the junction (along the direction labelled $x$ in Fig. 2a, inset), increasing (decreasing) when the projection is positive (negative). We assume that $n$ magnons impinge on the junction. We note that $p_i$ is the electric dipole moment of the $i$th magnon, $p_{\perp,i}$ is its projection along $x$ and we write $p_{\perp,i}^+$ (respectively $p_{\perp,i}^-$) when the electrical dipole moment projection along $x$ is positive (respectively negative). The sum of all the positive (negative) contributions is denoted as

$$I_{M,+} = \sum_{i=1}^{n^+} p_{\perp,i}^+,$$

$$I_{M,-} = \sum_{i=1}^{n^-} p_{\perp,i}^-,$$

where $I_{M,-} < 0$

where $n = n^+ + n^-$. This sum can be seen as an average positive (negative) dipole moment. The phase shift induced by magnons, denoted as $\delta \phi_{\text{M}}$, will then depend on the total electric dipole: $\delta \phi_{\text{M}} \propto |I_{M,+}| - |I_{M,-}|$; this corresponds to the shift shown in Fig. 2.

The strong visibility drop can, in turn, be understood by the fact that phase $\delta \phi_{\text{M}}$ fluctuates as magnons randomly impinge on the junction. Assuming a Poissonian emission of magnons, the fluctuations in $\delta \phi_{\text{M}}$ depend on the total amount of propagating magnons:

$$\delta \phi_{\text{M}} \propto |I_{M,+}| + |I_{M,-}|.$$  

where $|I_{M,+}|$ and $|I_{M,-}|$ are proportional to $|V_e| - E_p/e$. The oscillating term in the above transmission ($T_{\text{MZ}}$) then acquires an additional prefactor $e^{-i\phi_{\text{M}}}$ (ref. 19), yielding an exponential decrease in the MZI visibility as a function of $|V_e| - E_p/e$.

In Fig. 2f, we plot $\ln(V)$, where $V$ is the visibility shown in Fig. 2e, at $B \approx 8.875$ T. The perfect agreement with a linear dependence (Fig. 2f; red solid line) validates our assumption of a Poissonian source of magnons (note that using shot noise measurements, we have independently verified that the Poissonian assumption for the magnon emission is correct; see Supplementary Section X). The phase shift corresponding to these data, as shown in Fig. 2g, remains relatively small even above the threshold, implying that $|I_{M,+}| \approx |I_{M,-}|$. This can be understood by the fact that, on average, magnons perpendicularly impinge on the junction.

On the contrary, at a lower magnetic field ($B \approx 8.85$ T; Fig. 2b,c), the phase dependence with $V_e$ is much more pronounced. In this case, we compare the dependence of both $\ln(V)$ and phase shift as a function of $V_e$ (Fig. 2g), showing a perfect match between the two at our experimental accuracy. This match implies that $|I_{M,+}| + |I_{M,-}| \approx ||I_{M,+}| - |I_{M,-}||$, and signifies that magnons predominantly impinge on the junction from a given direction, which is determined later. The phase shift shown in Fig. 2b indicates that the magnetic field must be increased to keep the Aharonov–Bohm phase constant, implying that the area of the interferometer diminishes when magnons are emitted. Assuming a junction length of 1.5 μm, a lateral shift of one arm of the interferometer equals to 0.14 nm at $V_e = -2$ mV. This reduction in the interferometer area is similar to the effect of a positively biased electrostatic gate in the n-type region; thus, the projection of the electrical dipole moment of the impinging magnons along $x$ is mainly positive, that is, $I_{M,+} \gg I_{M,-}$, leading to $\delta \phi_{\text{M}} \propto I_{M,+}$.

In Fig. 2a and the inset, this is symbolized by the upper trajectory.

The comparison between Figs. 2c,e shows that, on one hand, the visibility only weakly depends on the magnetic field, implying that the total flux of magnons impinging on the junction remains roughly constant. On the other hand, the phase shift dramatically changes when modifying the field by 25 mT, indicating that the spatial distribution of the magnon trajectories is very sensitive to the magnetic field.
field. We have measured the phase shift for different points of the $\nu_b = 1$ plateau (Supplementary Figs. 17 and 18) and observed that a few tens of milliTeslas are enough to affect the magnon trajectory. Note that this phase shift can, thus, drastically change with the sign of $V_{BG}$, as this changes the emission site (as mentioned above) and thus the impinging magnons’ trajectory.

We now check that other dephasing mechanisms can be straightforwardly ruled out. One might naturally suspect that our observations are due to the absorption of magnons on a contact upstream of the interferometer, resulting in a finite d.c. bias $V_{m}$ injected into the system. A finite bias indeed is a source of phase shift and decoherence in MZIs\cite{15,24,25}

We can directly calculate $V_{m}$ by integrating the measured non-local differential voltage $\frac{dV_{xy}}{dV}$ across the MZI, where $V_{T}$ is the non-local voltage developing at the contact on which $T_{MZ}$ is measured and $V$ is the voltage applied at the magnon-emission contact (Fig. 2a). This yields the d.c. voltage developing at the output of the MZI, which we divide by the MZI transmission $T_{MZ}$ (measured independently) to obtain $V_{m} = \frac{2}{\tau_{T}} \int_{0}^{V_{T}} \frac{dV}{dV} dV$. At $V_{T} = 1.5 \text{mV}$, we find $V_{m} = 2 \mu \text{V}$.

Figure 3 shows that such a bias cannot be the source of the observed dephasing and decoherence. As depicted in Fig. 3a, we directly applied a finite d.c. bias $V_{T}$ at the input of the interferometer, and measured its oscillations as a function both $V_{T}$ and magnetic field. The oscillations, plotted in Fig. 3b, survive up to $V_{T} \approx 250 \mu \text{V}$, more than two orders of magnitude larger than the computed $V_{m}$.

This is further confirmed by computing the visibility (Fig. 3c), which shows the lobe structure typical of MZI over scales of 100 mV, as well as the phase (Fig. 3d), which also changes over voltage scales that are two orders of magnitude larger than the computed $V_{m}$.

On top of this finite d.c. bias $V_{m}$, there can also be fluctuations generated by the stochastic aspect of magnon absorption. The question whether fluctuations in the bias voltage applied to an MZI lead to additional decoherence has been previously studied, both theoretically\cite{25} and experimentally\cite{15,24,26}. It was shown that partitioning the edge channel upstream of the interferometer (and thus generating shot noise within it) did not affect the robustness of the interferences. Thus, if we assume that the fluctuations generated in our case are Poissonian (which are confirmed using noise measurements; Supplementary Section X), we can safely discard them as the source of our observed dephasing. Finally, our observations cannot be explained by local electromagnetic fluctuations caused by charge tunnelling at magnon emission and absorption. These fluctuations, as screened by the back gate and far from the interferometer, are essentially too small to drive decoherence (detailed calculations provided in Supplementary Section IX).

Here we discuss how the emission energy of magnons changes with the filling factor at around the $\nu_{b} = 1$ plateau. Figure 4a shows the Hall resistance $R_{H}$ around $\nu_{b} = 1$ as a function of the back-gate voltage $V_{BG}$. The MZI oscillations are measured for different $V_{BG}$ values (Fig. 4a, vertical dashed lines), and their corresponding visibility and phase data are summarized in Fig. 4d, e. We observe markedly different behaviours on the two edges of the $\nu_{b} = 1$ plateau. Figure 4b shows the oscillations measured close to $\nu_{b} = 2$ ($V_{BG} = -0.129 \text{V}$; Fig. 4a, blue vertical line), where the visibility starts dropping at a value of $V_{m}$ only slightly larger than that in the middle of the plateau (Figs. 2 and 4d) and the phase is weakly affected (Fig. 4e). At $V_{BG} = -0.18 \text{V}$, close to $\nu_{b} = 0$ (Fig. 4c and Fig. 4a, black vertical line), the visibility drops drastically at a much smaller $V_{m}$ value (<0.5 mV), and the phase shift is very pronounced (Fig. 4e).
The apparent decrease in the dephasing threshold as the electron density is tuned close to $\nu_B=0$ indicates that magnons with a gap smaller than the Zeeman energy are being excited. The cantile-antiferromagnetic phase of the $\nu=0$ quantum Hall state is generally expected to host gapless magnons $^{22-25}$ (additional details in Supplementary Section XII); the observation of the reduction in excitation gap at the transition from $\nu_B=1$ to $\nu_B=0$ is thus compatible with the transition towards such a phase. Note, however, that this approach does not directly probe the properties of $\nu=0$, since the electronic MZI is not defined at $\nu_B=0$ but rather the properties of the transient states between $\nu_B=0$ and $\nu_B=1$ that might already carry features of the $\nu_B=0$ state. Our results show that hallmark mesoscopic circuits such as electron interferometers provide new ways to probe the fundamental properties of quantum Hall ferromagnets.

**Online content**

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Methods
We used a CryoConcept dry dilution refrigerator with a base temperature of 13 mK. The measurements of transmitted currents and Hall resistance $R_H$ values were performed using multiple lock-in amplifiers with low-noise preamplifiers. Here a.c. excitations (1–10 nA) with different frequencies (70–300 Hz) were used. Buried ohmic contacts underneath the top gates enabled us to directly determine the filling factors from the regions of interest.

Data availability
Source data are provided with this paper. All the data, code and materials used in the analysis are available in some form to any researcher for purposes of reproducing or extending the analysis.

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Author contributions
A.A., M.J., P.B. and P. Rouleau performed the experiment with help from F.D.P. and P. Roche. A.A., P.B., M.J., F.D.P. and P. Rouleau analysed and discussed the data with help from P. Roche. T.J. and P. Rouleau developed the theoretical model. T.T. and K.W. provided the boron nitride layers. M.J. fabricated the device with inputs from P.B., A.A., F.D.P. and P. Rouleau. A.A., F.D.P., P. Roche and P. Rouleau wrote the manuscript with inputs from all the authors. P. Rouleau designed and supervised the project.

Competing interests
The authors declare no competing interests.

Additional information
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