Excitonic transport driven by repulsive dipolar interaction in a van der Waals heterostructure

Zhe Sun, Alberto Ciarrocchi, Fedele Tagarelli, Juan Francisco Gonzalez Marin, Kenji Watanabe, Takashi Taniguchi and Andras Kis

Dipolar bosonic gases are currently the focus of intensive research due to their interesting many-body physics in the quantum regime. Their experimental embodiments range from Rydberg atoms to GaAs double quantum wells and van der Waals heterostructures built from transition metal dichalcogenides. Although quantum gases are very dilute, mutual interactions between the particles could lead to exotic many-body phenomena such as Bose–Einstein condensation and high-temperature superfluidity. Here we report the effect of repulsive dipolar interactions on the dynamics of interlayer excitons in the dilute regime. By using spatially and temporally resolved photoluminescence imaging, we observe the dynamics of exciton transport, enabling a direct estimation of exciton mobility. The presence of interactions significantly modifies the diffusive transport of excitons, effectively acting as a source of drift force and enhancing the diffusion coefficient by one order of magnitude. Combining repulsive dipolar interactions with the electrical control of interlayer excitons opens up appealing new perspectives for excitonic devices.

Van der Waals (vdW) heterostructures with a type-II band alignment enable the formation of long-lived interlayer excitons (IXs) composed of charges that are spatially separated in distinct layers. In contrast to intralayer excitons in monolayer transition metal dichalcogenides, the spatial separation of charges gives rise to a sizable permanent out-of-plane electrical dipole moment, which makes IXs a promising platform to realize electrically controlled excitonic devices. These are a new class of solid-state devices for information and signal processing, analogous to electronic or spintronic devices, but based on encoding information in the amplitude and/or pseudo-spin of the exciton currents that can be controlled using electrical fields. As a result, exciton transport in vdW heterostructures has recently attracted growing interest. Many basic questions, however, remain open, such as the nature of mutual interactions between IXs. Moreover, exciton transport in vdW heterostructures has so far been attributed only to exciton diffusion currents and the observation of long IX transport distances is usually assigned to a large effective diffusion coefficient—a phenomenological parameter that conceals the role of the repulsive dipolar interaction on IX transport. In analogy with semiconducting devices, it would also be advantageous to introduce exciton drift currents to increase the range of exciton motion in a device and enhance control over exciton transport.

Here, by imaging the temporal evolution of the IX cloud, we reveal repulsive dipolar exciton–exciton interactions as the driving force behind IX transport, acting as an effective source of a drift field. Concurrently, we directly deduce the exciton mobility from power-dependent drift velocities. Our findings, combined with the electrical control of IXs, pave the way towards the control of the motion of IXs over long distances.

In our work, instead of investigating IXs in a heterobilayer, we introduced monolayer hexagonal boron nitride (1L-hBN) as a thin spacer between the monolayers of WSe₂ (1L-MoSe₂) and monolayers of MoSe₂ (1L-MoSe₂). The motivations are twofold: first, the spacer can increase the separation between electrons and holes, enhancing the size of the electrical dipole by a factor of ~1.5 (ref. 5). Second, the presence of a moiré potential with an amplitude of the order of 100–200 meV (refs. 9–11) in MoSe₂/WSe₂ heterostructures localizes IXs by effectively decreasing the diffusion coefficient.¹ The Mois weakens the moiré potential and reduces its period due to the large lattice mismatch between hBN and MoSe₂/WSe₂ while retaining a sufficiently strong transition dipole moment for hosting bright IXs.

Figure 1a,b shows an optical image and schematic of the device. It consists of a WSe₂/hBN/MoSe₂ heterotrilayer encapsulated in hBN, with a transparent global top gate and several local back gates. Multiple local back gates allow us to control the exciton flux by applying a laterally modulated vertical electric field (Fig. 1b). The yellow-shaded area in Fig. 1a indicates the heterotrilayer region where we perform the spatially and temporally resolved photoluminescence (PL) measurements at 4.6 K. We use a sub-picosecond 725 nm pulsed laser with an 80 MHz repetition rate to excite the left part of the heterostructure and to generate an initial population of IXs (Fig. 1a, red spot). Figure 1c presents the photoluminescence (PL) intensity of IXs at the excitation spot as a function of the emission wavelength and average laser power \( P_{\text{ave}} \). In Fig. 1d, we show the peak intensity and peak energy extracted from the data in Fig. 1c. With increasing laser power, the peak intensity first increases linearly and then begins to saturate at higher laser powers (>300 \( \mu \)W).

The observed blueshift in the IX energy with increasing laser power is due to the repulsive exciton–exciton interaction, which can be decomposed into two terms: dipolar repulsion (which is valley independent) and exchange interaction (which is determined by the valley indices)¹⁰. The dipolar interaction is purely repulsive and can be estimated using a parallel-plate capacitor model. This gives a lower bound on the exciton density, as it does not account...
for a reduction in interaction energy due to a rearrangement of the interlayer excitons caused by Coulomb repulsion. Following the rearrangement and reduction in interaction energy, a larger density of IXs will be required to achieve the same energy shift. The exchange interaction has a more complex dependence on the electron–hole separation since it can change from repulsive to attractive. As the vertical separation of electrons and holes becomes larger, the exchange interaction decreases and becomes negative when the vertical separation is larger than the Bohr radius of IX, namely, $a_0 \approx 1\text{ nm}$ (refs. 19,20). In our case, since the separation between 1L-WSe$_2$ and 1L-MoSe$_2$ is around 0.9 nm and similar to the Bohr radius of IX, we can neglect the exchange interaction and only consider the dipole–dipole interaction.

To quantify the influence of dipole–dipole interactions on the exciton transport using time-resolved imaging, we first estimate the exciton initial density $n_0$ from the blueshift $\delta E_{XX}$ via the parallel-plate capacitor model$^{17–19}$:

$$\delta E_{XX} = n_0 U_{XX} = n_0 \frac{e^2 d}{\varepsilon_0 \varepsilon_{HS}},$$

where $e$ is the elementary charge of an electron, $d=0.9$ nm is the out-of-plane dipole size of IXs, $\varepsilon_0$ is the vacuum permittivity, $\varepsilon_{HS}=6.26$ is the effective relative permittivity of the WSe$_2$/hBN/MoSe$_2$ heterotrilayer and $U_{XX} \approx 2.6\mu\text{eV/}\mu\text{m}^2$ is the exciton–exciton interaction strength (Supplementary Note 1). At $P_{\text{ave}} = 50, 100$...
and 200 μW, we determine the exciton densities as about 2, 4 and $8 \times 10^{11}$ cm$^{-2}$, respectively. The exciton densities extracted from the spectral shift are consistent with the values estimated from the applied laser powers (Supplementary Note 2).

To image the spatial and temporal distribution of IXs, we use the setup depicted in Supplementary Fig. 3, in which the emitted photons are filtered ($< 1.45$ eV) and sent to either a charge-coupled device (CCD) camera or a home-made scanning avalanche photodiode (APD) system (Supplementary Note 3). Figure 1fg shows the CCD images of the normalized PL emission intensity from IXs, acquired for different excitation powers. Compared with the CCD image of the focused excitation spot (Fig. 1e), the spatial profile of IXs extends farther and exhibits a growing size with increasing excitation power, signalling the presence of strong repulsive exciton–exciton interactions. A comparison between the power-dependent PL spectrum for MoSe$_2$ intralayer excitons and IXs in the heterolayer region also proves the existence of strong interactions between the IXs (Supplementary Note 4).

To acquire a map of the PL intensity as a function of position and time $I(x,y,t)$, the APD is scanned in the image plane across the emission spot$^{12,23}$. We use a time-correlated photon-counting module to record the photon clicks. To rule out the decay of PL intensity induced by radiative emission, the raw data $I(x,y,t)$, proportional to the exciton density distribution $n(x,y,t)$, are normalized at each recorded time $t_i$ to obtain $I_{norm}(x,y,t_i) = I(x,y,t_i)/\max(I(x,y,t_i))$ (Supplementary Note 5). Figure 2a shows the two-dimensional (2D) spatial profiles of IXs at different times and $P_{in} = 200$ μW. To further analyse the expansion of the spatial profile, the area of IXs at each time is extracted by counting the number of pixels for which the normalized PL intensity $I_{norm}(x,y,t)$ is higher than 0.2. This is shown in Fig. 2c for three different excitation powers. In the absence of spatial constraints$^{12,23}$, the exciton cloud is expected to grow linearly as a function of time, with the evolution of exciton density $n(x,y,t)$ due to the diffusion described by

$$\frac{\partial n}{\partial t} = D \nabla^2 n - \frac{n}{\tau},$$

where $D$ denotes the diffusion coefficient and $\tau$ is the exciton radiative lifetime.

We find, however, that the area occupied by the IX cloud initially ($t < 2$ ns) grows at a higher, power-dependent speed, but then slows down to a speed that is independent of the excitation power and initial exciton density. Even though, in our case, the motion of IXs is limited by the finite size of the heterostructure, this cannot explain the significant power dependency at early times, when the expansion is not constricted by the edges. Instead, we attribute the observed exciton cloud dynamics to dipolar interactions.

A sublinear increase in the IX cloud area with time has been observed before$^{12,25}$ and has been attributed to two possible mechanisms. First is the effect of the strong moiré potential introducing a modification of diffusivity

$$D = D_0 e^{-\frac{U_{moir}}{U_{max} T^\frac{1}{2}}},$$

where $D_0$ is the Boltzmann constant, $T$ is the temperature, $D_0$ denotes the bare diffusivity and $U_{moir}$ is the depth of the moiré trapping potential$^{14}$. The second is the generation of electron–hole plasma at exciton densities exceeding the Mott transition density $n_{Mott} \approx 10^{13}$ cm$^{-2}$ (refs. 25–28). Neither of these reports, however, show a large area of transport ($> 5$ μm$^2$) in dilute excitonic gases ($n_e \leq n_{Mott}$). We emphasize that in our work, in contrast with previous results, we observe a distinct power-dependent expansion of the IX cloud in the dilute regime. In addition, in the heterolayer, the 1L-HBN spacer between WSe$_2$ and MoSe$_2$ is expected to weaken trapping due to the moiré potential by increasing the spatial separation between the electron and hole wavefunctions, thereby facilitating the propagation of IXs$^1$. Numerical simulations that include the effect of the moiré potential fail to reproduce our data (Supplementary Note 6). We, therefore, explain the exciton transport only using repulsive dipolar interactions, decreasing in strength as the IXs cloud expands and radiatively decays.

To identify the contribution of exciton–exciton interactions, we introduce a power-related term into a 2D drift–diffusion equation and solve it numerically (Supplementary Note 7). The drift–diffusion equation describes the spatial and temporal distribution of exciton density $n(x,y,t)$:

$$\frac{\partial n}{\partial t} = D \nabla^2 n + \frac{\mu}{e} \nabla (n \nabla \delta E) - \frac{n}{\tau},$$

where $\mu$ is the exciton mobility that can be expressed using $D$ and temperature $T$ via the Einstein relation $\mu = De/k_b T$ and $\delta E$ is the total potential energy of IXs. The first term on the right-hand side of equation (2) is the diffusion term, whereas the second term denotes the drift term. Here $\delta E = \delta E_{XX} = n(x,y,t)U_{XX}$, leading to an exciton potential energy that varies in both time and space. The excitation power enters the equation via the initial exciton density $n_0(x,y)$, which is expected to be of the same order of magnitude as the exciton density determined from the interaction-induced energy shift (Fig. 1d). We determine the lifetime of IXs to be about $t = 3.5$ ns from the time-resolved PL measurements. Since the IX lifetimes in the studied power range do not decrease with an increase in exciton density, we neglect the effect of exciton–exciton annihilation$^{11}$ (Supplementary Note 8).

To highlight the effect of the drift term on exciton transport, Fig. 2b shows the simulated exciton area as a function of time with and without $\delta E_{XX}$ (solid and dashed lines, respectively). We use $U_{XX} = 2.6 \mu$eV μm$^2$ and $n_0 = 1 \times 10^{11}$ cm$^{-2}$ and do not take the finite size of the heterolayer into consideration. When $U_{XX}$ is neglected, the exciton area increases linearly with time and the slope is proportional to $D$. With $U_{XX}$ included, the area first increases sublinearly; however, at later times ($> 2.5$ ns), when the drift term almost vanishes, the slope becomes the same as in the case of neglected $U_{XX}$. Similar simulations have successfully reproduced the experimental results of indirect exciton diffusion in GaAs double quantum wells with strong dipole–dipole interactions$^{27}$. We fit our data in Fig. 2c using the same model and take the boundaries of the heterostructure into account (Supplementary Note 9). To be consistent, we employ the same condition ($n_{norm}(x,y,t = t_i) > 0.2$) to calculate the exciton area. In Fig. 2c, we use $D = 0.15$ cm$^2$ s$^{-1}$ and $U_{XX} = 2.6 \mu$eV μm$^2$ as the parameters and treat the initial exciton density $n_0$ as a free variable to fit our data. The deduced exciton densities are similar to the values determined from the blueshift energy and the applied laser power, confirming the consistency of our model. Considering the effective diffusion coefficient, this is enhanced by a factor of ~12 for $P_{in} = 200$ μW when $t \leq 1$ ns (the ratio of the slope between the red dashed line and black dashed line).

We use the Einstein relation first to estimate the exciton mobility from the diffusion constant, finding $\mu \approx 380$ cm$^2$ (V s)$^{-1}$, similar to the low-temperature mobility of single charge carriers in monolayer transition metal dichalcogenides$^{29–30}$. This indicates that at this temperature, exciton transport could be limited by the same mechanism, namely, charged-impurity scattering.

We now turn to the drift term and quantify the relationships between the exciton drift velocities and mobility, in analogy with the transport of charge carriers in semiconductors. To visualize this relationship, we use an alternative way to normalize the raw data $I(x,y,t)$ by normalizing each spatial coordinate $(x,y)$ by $I_{norm}(x,y,t) = I(x,y,t)/\max(I(x_0,y,t))$ (Supplementary Note 5). This allows us to observe the spatial distribution of the IX lifetime as well as the spatially and temporally resolved IX transport by
scanning the APD along \( x = 0 \) in Fig. 2a with a finer step, we obtain the one-dimensional (1D) normalized PL data \( f_{\text{time}}(0, y, t) \) and \( f_{\text{space}}(0, y, t) \) for different excitation powers, as shown in Fig. 3a–c. From \( f_{\text{time}} \), we observe that IXs propagate to about \( y_{\text{drift}} \approx 1.5 \mu m \) in the \(+y\) direction, as indicated by the white dashed lines in Fig. 3a–c; the bottom panels distinctly show the time delay at \( y > 0 \) generated during IX propagation in the \(+y\) direction. Figure 3d shows the simulated normalized exciton density distribution \( f_{\text{time}}(0, y, t) \) and \( n_{\text{space}}(0, y, t) \) for \( n_0 = 4 \times 10^{11} \text{ cm}^{-2} \) using equation (2). Further, \( f_{\text{space}} \) allows us to extract the effective speed \( v_{\text{eff}} \) of IXs at different powers (Supplementary Note 10). The black dashed lines in the bottom panels of Fig. 3a–c are guides for the eye, which highlight the effect of excitation power and initial exciton density on \( v_{\text{eff}} \). We can further decompose \( v_{\text{eff}} \) into its diffusion and drift components as \( v_{\text{eff}} = v_{\text{diff}} + v_{\text{drift}} \). At a higher excitation power, due to repulsive interactions, IXs experience a stronger effective drift field \( F_{\text{drift}} \), which can be deduced from the spectral blueshift \( \delta E_{\text{xx}} \) and drift distance \( y_{\text{drift}} \) via \( F_{\text{drift}} = \delta E_{\text{xx}}/y_{\text{drift}} \). In analogy with the transport of charge carriers in an electric field, the neutral exciton mobility \( \mu \) can be approximately expressed using \( \mu = v_{\text{eff}}/F_{\text{drift}} \), allowing us to directly estimate the exciton mobility without using the Einstein relation.

We extract \( v_{\text{eff}} \) at different powers from \( f_{\text{space}}(0, y, t) \) and plot them with the spectral blueshift \( \delta E_{\text{xx}} \) in Fig. 3e. The error bars of \( v_{\text{eff}} \) are given by the linear fits along the black dashed lines. By applying a linear fit to the data in Fig. 3e, we estimate \( \mu \approx 440 \text{ cm}^2 \text{s}^{-1} \text{V}^{-1} \text{s}^{-1} \), which is consistent with the value calculated from the diffusion coefficient. The presence of the power-dependent excitonic drift force shows that the repulsive dipolar interactions can be used to control IX transport.

Next, we present how combining the repulsive interactions together with an external electric field can be used to control the motion of IXs. We generate a modulated electrostatic potential \( \delta E_{\text{oe}}(x) \) along the \( x \) direction using one of the local back gates, which creates a spatially varied but time-independent energy profile acting...
Fig. 3 | Drift velocity and mobility of interlayer excitons. a–c. One-dimensional normalized PL intensity along \( x = 0 \) in Fig. 2a for \( P_{\text{ex}} = 100, 200 \text{ and } 350 \mu\text{W} \). Top: \( n_{\text{norm}} \); bottom: \( n_{\text{form}} \). The white dashed lines enclose the region of IX transport. The black dashed lines are guides for the eye for \( \nu_{\text{eff}} \). Simulation of 1D normalized exciton distribution along \( x = \delta \) using equation (2) for \( n_{\text{ex}} = 4 \times 10^{10} \text{ cm}^{-2} \). Top: \( n_{\text{form}} \); bottom: \( n_{\text{norm}} \). d. Effective exciton velocity \( \nu_{\text{eff}} \) as a function of the spectral blueshift \( \delta E_{\text{XX}} \) extracted from Fig. 1d. The error bars of \( \nu_{\text{eff}} \) are given by the linear fits along the black dashed lines in a–c (Supplementary Fig. 12). The red solid line is a linear fit to the data.

on the IXs: \( \delta E_{\text{el}}(x) = -edF_{\text{el}}(x) \), where \( F_{\text{el}} \) is the external electric field (ref. 1). Figure 4a,b, top, shows the schematic of the energy profiles as well as the expected exciton motion for back-gate voltages of \( V_{\text{bg}} = -2 \) and 2 V, respectively. Figure 4a,b, middle, presents the CCD images of the normalized IX PL intensity using \( P_{\text{ex}} = 200 \mu\text{W} \). The region enclosed by the black dashed lines indicates the position of the local back gate. By tuning the gate region higher or lower in energy with respect to its surroundings, we generate a potential barrier or a trap, effectively controlling the propagation distance of IXs along the \( +x \) direction. We measure the 1D normalized PL intensity \( n_{\text{form}}(x, 0, t) \) along \( y = 0 \) using the scanning APD system. As shown in Fig. 4a,b, bottom, we clearly observe the process of IX gas moving into the back-gate region as we adjust the electrostatic potential configuration from a barrier to a trap. Instead, Fig. 4c shows the 1D normalized PL intensity \( n_{\text{norm}}(x, 0, t) \) for \( V_{\text{bg}} = -1 \text{ V} \), which is a barrier configuration of lower amplitude with respect to the data in Fig. 4a. Here the excitons first flow towards the gate region of lower potential, but then surprisingly move away from the gate region after \( \sim 1 \text{ ns} \). This phenomenon is another manifestation of the repulsive exciton–exciton interaction, and its interplay with the electrostatic potential in the device. To explain this observation, in Fig. 4e, we show a schematic of the energy profile of dipolar repulsion \( \delta E_{\text{XX}}(t) = n(t)U_{\text{XX}} \) and electrostatic energy \( \delta E_{\text{el}} \). Initially, when the exciton density is high, \( \delta E_{\text{XX}} \) dominates over \( \delta E_{\text{el}} \), such that the exciton flux is pushed towards the gate region. As the exciton density decreases due to spatial expansion and radiative emission, \( \delta E_{\text{el}} \) becomes higher than \( \delta E_{\text{XX}} \), resulting in exciton flux leaving the gate region. The observation of this competitive phenomenon requires \( n_{\text{ex}}U_{\text{XX}} \) to be larger than but comparable to \( \delta E_{\text{el}} \). In this measurement, we create such a condition by using \( P_{\text{ex}} = 200 \mu\text{W} \), which corresponds to \( \delta E_{\text{XX}}(t = 0) = \delta E_{\text{el}} \approx 20 \text{ meV} \) and \( \delta E_{\text{el}} \approx 9 \text{ meV} \) for \( V_{\text{bg}} = -1 \text{ V} \) (Supplementary Note 1). We introduce \( \delta E_{\text{el}} \) into the simulation to better evaluate the interplay between \( \delta E_{\text{XX}} \) and \( \delta E_{\text{el}} \). Here we assume that the electrostatic potential has the form of a harmonic potential \( \delta E_{\text{el}}(x - \chi)^2 \) with height \( \delta E_{\text{el}} \) (Fig. 4f, inset). \( \chi \) is the coordinate of the potential center. Figure 4f shows the experi-
mental results as well as simulations of the propagation distance $L_{\text{es}}$ along the $+x$ direction as a function of $\delta E_{\text{el}}/\delta E_{\text{XX}}$ (Supplementary Note 11); therefore, adjusting this ratio allows us to precisely control the propagation distance of IXs.

Spatially and temporally resolved PL imaging performed in the heterobilayer region of the same sample (Supplementary Note 12) shows no exciton diffusion. Under low excitation intensities and in contrast to the spectra from the heterotrilayer region, we find multiple narrow peaks, with linewidths of 0.5–1.0 meV, limited by the energy resolution of our spectrometer grating. The narrow peaks are sensitive to electrostatic doping and could originate from excitons localized by strain or moiré traps. Indeed, moiré trions have been recently observed by several groups in WS$_2$/MoSe$_2$ heterobilayers$^{32,33}$. Exciton trapping explains the absence of a sizable expansion of the IX cloud in the heterobilayer region. Moreover, this further demonstrates that the 1L-hBN spacer between the WS$_2$ and MoSe$_2$ layers could strongly reduce the effect of moiré potential and strain traps, resulting in an evident excitonic diffusion/drift.

Our results demonstrate that repulsive dipolar interactions in dilute excitonic gases have a strong influence on exciton transport and that they act as a source of drift force. Time-resolved PL imaging enables us to visualize the dynamic evolution of IXs in the WS$_2$/hBN/MoSe$_2$ heterotrilayer and allows us to quantify the diffusion coefficient and exciton mobility that play a central role in the prospect for applications of excitonic devices. Our findings constitute a crucial step towards the use of spatial patterns of a laser field to control the propagation of IXs$^{34,35}$, for example, by using doughnut-shaped beams as simulated and discussed in more detail in Supplementary Note 13. The excitons that are driven away from the laser hotspot by repulsive exciton-exciton interactions are expected to constitute a cold excitonic gas. Many exotic phases of matter and emergent phenomena might appear in it, including Bose–Einstein condensation and high-temperature superfluidity$^{36–38}$. Finally, a strong spatial confinement of IXs that experience strong repulsive dipolar interactions leads to nonlinearities in energy$^9$. Due to the additional 1L-hBN spacer and thus stronger repulsive interactions, we expect the nonlinearity of localized excitons in the heterotrilayer to be more pronounced.

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Methods

Device fabrication. Thin Cr/Pt (2–3 nm) local back gates were patterned using electron-beam lithography and deposited on a silicon substrate using electron-beam evaporation. The heterostructure was fabricated using a polymer-assisted transfer method. Flakes were first exfoliated on a polymer double layer. After the monolayers were optically identified by PL measurements, the bottom layer was dissolved with a solvent and free-floating films with flakes were obtained. These were transferred using a home-made setup with micromanipulators to carefully align the flakes on top of each other. The polymer residue was removed with a hot acetone bath. Afterwards, the heterostructure was thermally annealed under high-vacuum conditions ($10^{-6}$ mbar) for 6 h at 340 °C. Finally, electrical contacts (80 nm Pd for contacts to the flakes and 8 nm Pt for the global top gate) were patterned using electron-beam lithography and deposited using electron-beam evaporation.

Time-resolved optical characterization. A confocal microscope is used to optically excite IXs and collect the emitted photons through the same objective with a working distance of 4.5 mm and numerical aperture of 0.65. IXs are excited using sub-picosecond pulses with a repetition rate of 80 MHz from a Ti:sapphire laser (Coherent, Chameleon Ultra). The collected photons are sent to an APD (Excelitas Technologies, SPCM-AQRH-16) mounted on a 2D motorized translational stage. The output of the APD is connected to a time-correlated photon-counting module with a resolution of 12 ps r.m.s. (PicoQuant, PicoHarp 300), which measures the arrival time of each photon. For the measurements in this work, we set the time bin to 64 ps to record the photon clicks. The single-photon timing resolution of the APD is ~350 ps, which is the main time limitation for this setup. The technical details can be found in Supplementary Note 3.

Numerical simulations. The 2D drift–diffusion equations are numerically solved using the forward-time centred-space method. This method is based on the central difference in space and the forward Euler method in time. The relevant details can be found in Supplementary Note 6.

Data availability

The data that support the findings of this study are available from the corresponding authors upon reasonable request.

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Author contributions

Z.S. built the experimental setups, performed the optical measurements and analysed the data with input from A.K. A.K. initiated and supervised the project. A.C. fabricated the device. F.T. worked on device fabrication. J.F.G.M. contributed to the initial stages of setup development. K.W. and T.T. grew the hBN crystals. Z.S. and A.K. wrote the manuscript with input from all the authors.

Competing interests

The authors declare no competing interests.

Additional information

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Correspondence and requests for materials should be addressed to Zhe Sun or Andras Kis.

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