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# Photonic indistinguishability of the tin-vacancy center in nanostructured diamond

Jesús Arjona Martínez<sup>1,\*</sup>, Ryan A. Parker<sup>1,\*</sup>, Kevin C. Chen<sup>2</sup>, Carola M. Purser<sup>1</sup>, Linsen Li<sup>2</sup>, Cathryn P. Michaels<sup>1</sup>, Alexander M. Stramma<sup>1</sup>, Romain Debroux<sup>1</sup>, Isaac B. Harris<sup>2</sup>, Martin Hayhurst Appel<sup>1</sup>, Eleanor C. Nichols<sup>1</sup>, Matthew E. Trusheim<sup>2</sup>, Dorian A. Gangloff<sup>1,3,†</sup>, Dirk Englund<sup>2,†</sup>, and Mete Atatüre<sup>1,†</sup>

<sup>1</sup>*Cavendish Laboratory, University of Cambridge,*

*JJ Thomson Avenue, Cambridge CB3 0HE, United Kingdom*

<sup>2</sup>*Department of Electrical Engineering and Computer Science,*

*Massachusetts Institute of Technology, Cambridge, Massachusetts 02139, USA and*

<sup>3</sup>*Department of Engineering Science, University of Oxford, Parks Road, Oxford, OX1 3PJ*

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Tin-vacancy centers in diamond are promising spin-photon interfaces owing to their high quantum-efficiency, large Debye-Waller factor, and compatibility with photonic nanostructuring. Benchmarking their single-photon indistinguishability is a key challenge for future applications. Here, we report the generation of single photons with  $99.7_{-2.5}^{+0.3}\%$  purity and 63(9)% indistinguishability from a resonantly excited tin-vacancy center in a single-mode waveguide. We obtain quantum control of the optical transition with 1.71(1)-ns-long  $\pi$ -pulses of 77.1(8)% fidelity and show it is spectrally stable over 100 ms. A modest Purcell enhancement factor of 12 would enhance the indistinguishability to 95%.

Indistinguishable photons from quantum emitters provide a fundamental resource for scalable quantum communication and have been employed to realize linear optical quantum computation [1–3], spin-photon and spin-spin entanglement [4–6], and quantum repeater schemes [7–10]. Experimentally, the photonic indistinguishability can be benchmarked through two-photon quantum interference known as the Hong-Ou-Mandel (HOM) effect. The HOM indistinguishability places a bound on the fidelities achievable in photon-mediated gates and entanglement distribution in measurement-based protocols [11–13]. This effect has been observed across multiple solid-state emitters such as the nitrogen-, silicon- and germanium-vacancies in diamond [14–16], defects in silicon-carbide [17], and semiconductor quantum dots [18].

Within solid-state emitters, the negatively charged group-IV centers in diamond stand as promising spin-photon interfaces due to their large Debye-Waller factor (60% - 80%) [19–22], competitive quantum efficiency (10% - 80%) [21–24] and first-order insensitivity to electric-field noise [25, 26]. This electric-field agnosticism is compatible with complex photonic nanostructuring, as charge noise does not couple deleteriously to a proximate emitter. Accordingly, high collection efficiency [27, 28], Purcell enhancement [29, 30], and incorporation into photonic integrated circuits [31] have been demonstrated using group-IV centers embedded in diamond. The negatively charged tin-vacancy (SnV) is particularly promising. Its large ground-state orbital splitting inhibits phonon-mediated dephasing [32], allowing for operation at temperatures accessible in standard helium closed-cycle cryostats. Accordingly, it has shown a spin-coherence time ( $T_2$ ) of 0.30(8) ms at 1.7K [33]; outperforming other group-IV centers at this temperature. Observation of transform-limited emission in nanopillars [34] and advancements in fabrication and charge stability [35] further demonstrate that SnV centers are a suit-

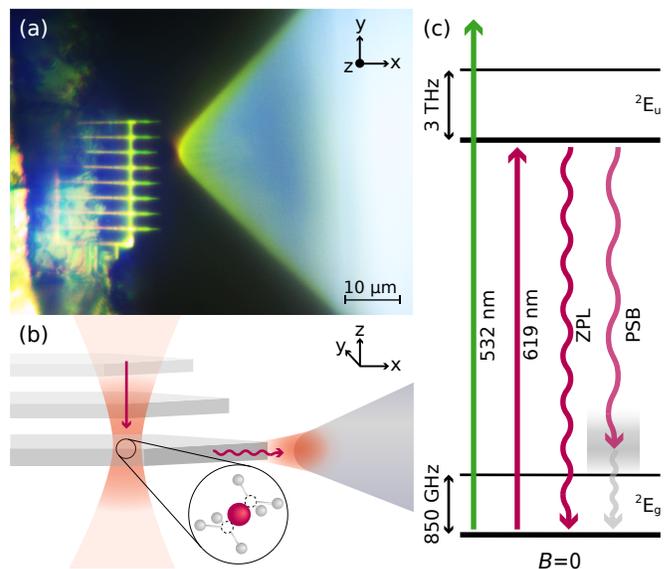


FIG. 1. (a) Microscope photograph of the lensed fiber aligned to a single waveguide using a 3-axis nanopositioning stack. (b) Diagram of the excitation and collection geometry. (c) Electronic structure of the SnV center with no magnetic field. Off-resonant 532-nm and on-resonance 619-nm lasers generate emission into the ZPL and PSB radiative decay pathways.

able spin-photon interface for quantum networking and measurement-based computation [12]. Confirming the optical coherence of the emitted photons is a necessary step towards these goals.

In this work, we report the observation of quantum interference of single photons from a SnV center in a single-mode diamond waveguide with an indistinguishability of 63(9)% and a single-photon purity of  $99.7_{-2.5}^{+0.3}\%$ . The orthogonal propagation directions of the excitation and collection modes, enabled by the waveguide nanostructure, suppress the excitation light by over 60dB. We further realize coherent control of the optical transition

of the SnV center with a  $\pi$ -rotation fidelity of 77.1(8)% performed in 1.71(1) ns. Thus, we show that the SnV center has sufficient photonic coherence to satisfy the requirements for quantum networking.

We perform our experiments on a diamond grown via chemical vapor deposition, implanted with Sn<sup>++</sup> ions, annealed, and fabricated into a waveguide chiplet [36, 37]. A chiplet consists of eight single-mode waveguides, each 50  $\mu\text{m}$  long with a rectangular cross-section of 200 nm by 270 nm. Figure 1(a) shows a characteristic device. The waveguides taper adiabatically over 9  $\mu\text{m}$  at the ends to maximize the theoretical out-coupling efficiency [37]. A support diamond structure connects the waveguides, allowing the chiplet to be pick-and-placed onto a silicon substrate edge [31]. The SnV center fluorescence couples into a single-mode, lensed fiber giving a 23(3)-fold enhancement in collection efficiency relative to conventional confocal microscopy performed on the same emitter studied in this work [37].

Figure 1(b) shows the orthogonal excitation and collection directions. In this geometry, the continuous-wave laser suppression is over 60 dB. The fiber-coupled SnV fluorescence is then routed from inside the cryostat, held at 3.6 K, to our optical setup [37].

In the absence of magnetic field, the SnV center has a spin-degenerate optical transition between the lower orbital branch of the ground state and the lower orbital branch of the excited state [52]. A 619-nm laser excites this transition resonantly, as highlighted in Fig. 1(c). Also highlighted are the two radiative-decay pathways, the zero-phonon line (ZPL) and phonon sideband (PSB). Resonant drive induces occasional blinking. To remedy this we alternate between resonant and off-resonant 532-nm excitation [37], which pumps the emitter into the photo-active, -1 charge state [35]. The studied emitter has an excited-state lifetime  $T_1$  of 7.44(20) ns, which corresponds to a transform-limited linewidth  $\Gamma_0/2\pi$  of 21.4(2) MHz [37]. This is longer than the 5.0 ns reported for bulk diamond [22], likely due to the emitter being positioned close to the diamond-air interface [22, 53].

We first demonstrate optical control of the SnV center through resonant excitation and collection of the PSB. Figure 2(a) shows the averaged PSB fluorescence during a 20-ns-long resonant 619-nm pulse starting at time  $\tau = 0$ . The histogrammed time-resolved fluorescence throughout the pulse is proportional to the instantaneous excited-state population and shows Rabi oscillations between the ground and excited states. The laser power  $P$  is parameterized by the saturation parameter  $s = P/P_{\text{sat}}$ , where  $P_{\text{sat}} = 31(2)$  nW is the resonant laser power at which the Rabi rate  $\Omega$  equals  $\Gamma_0/\sqrt{2}$ . For each power the excited-state population is fit by a two-level master equation, to extract the corresponding  $\Omega$  and dephasing rate [37, 54, 55]. An optical  $\pi$ -rotation at the highest driving power ( $s = 367(1)$ ), is performed in 1.71(1) ns, much faster than  $T_1$ , with a fidelity of 77.1(8)%.

To investigate the effects of decoherence, the master equation of Fig. 2(a) includes spontaneous emission and

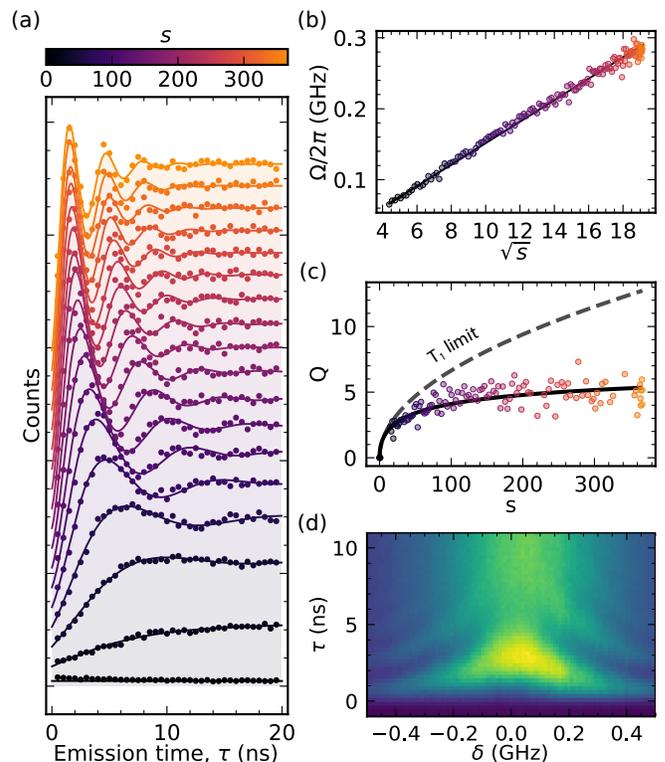


FIG. 2. (a) Histogrammed PSB fluorescence during a 20-ns-long resonant laser pulse. Curves are offset proportionally to the saturation parameter,  $s$ . Solid curves are fits to a master equation [37]. At the highest saturation parameter the population saturates to 0.5 at long times [37]. (b) Optical Rabi frequency  $\Omega$  extracted from panel a as a function of  $\sqrt{s}$ . The black line is a linear fit with zero intercept. (c) Quality factor  $Q$  as a function of  $s$ . The solid black curve is a fit to a master-equation including a pure dephasing rate proportional to  $\Omega$ . The dashed black curve shows the absolute coherence limit. (d) Detuning  $\delta$  dependence of the Rabi oscillations at  $s = 102(4)$  for the same excitation scheme as panel a. Color-coding with blue (yellow) corresponds to low (high) fluorescence.

pure dephasing, as well as a shot-to-shot detuning fluctuation resulting in inhomogeneous dephasing [37]. We confirm in Fig. 2(b) a direct proportionality of the Rabi rate on the square root of the saturation parameter through the relationship  $\Omega = \Gamma_0\sqrt{s}/2$ . This linear relationship extends to our highest probed  $s$  and indicates that control-limiting imperfections such as phonon-coupling or multi-level driving do not cause an appreciable deviation from the master-equation model [37]. Figure 2(c) shows the power dependence of the quality factor  $Q$ , defined as the product of  $\Omega$  and the  $1/e$  envelope decay time extracted from the master-equation fits. In the low excitation power regime,  $Q$  increases with power and agrees with the limit set by  $T_1$ . At higher powers  $Q$  saturates, which implies a laser-induced dephasing mechanism best modeled with a rate that linearly depends on  $\Omega$ . Possible mechanisms could include photo-induced charge noise or resonant coupling to phononic modes [56–58].

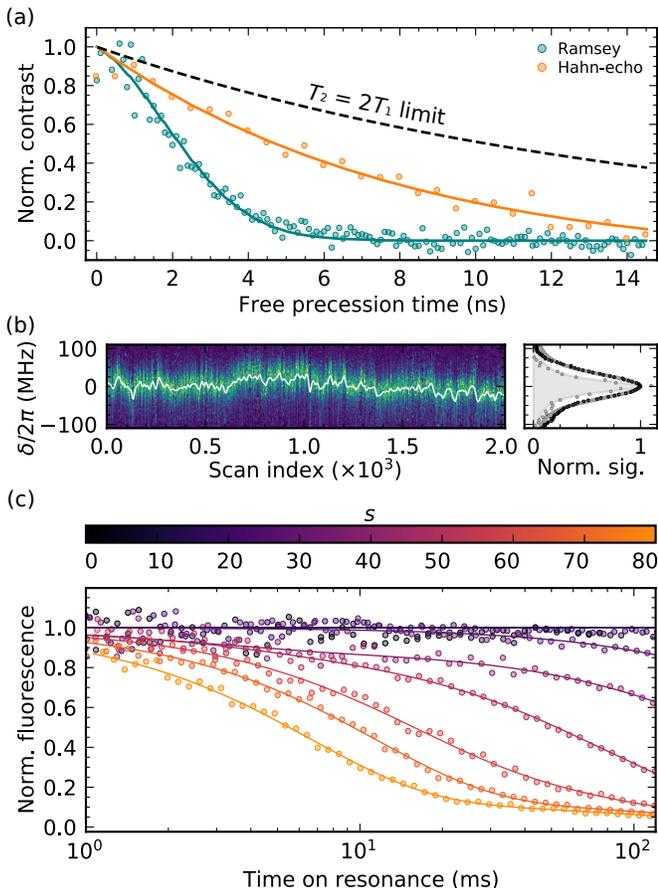


FIG. 3. (a) Measured Ramsey (teal) and Hahn-echo (orange) contrast decay envelopes with master equation fits (solid curves) [37]. The dashed black curve is the absolute coherence limit,  $T_2 = 2T_1$ . The contrast at 0 ns is normalized to one. (b) Left: Shot-to-shot evolution of the PLE lineshape. Each vertical cut is the average of fast PLE scans over 500 ms. Off-resonant 532-nm excitation is applied between line-cuts. The solid white line follows the emitter resonance frequency. Right: Example line-cut (gray) and inhomogeneous distribution of all line-cuts (black), with Lorentzian and Gaussian fits respectively. (c) Average PLE intensity as a function of time following a 532-nm off-resonant pulse for multiple saturation parameters. Solid curves are fits to bi-exponential decays. The experimental protocol is identical to that of panel b. The dark-state pumping time is normalized by accounting for the fraction of time spent on resonance.

By varying the frequency of the resonant laser, we also probe the dependence of the fluorescence on the detuning  $\delta = \omega_1 - \omega_0$ , where  $\omega_1$  and  $\omega_0$  are the angular frequencies of the 619-nm resonant laser and the transition, respectively. Figure 2(d) shows the time evolution of the fluorescence as a function of  $\delta$  at  $s = 102(4)$ , yielding a Rabi rate of  $153(3)$  MHz at  $\delta = 0$  and faster, lower-amplitude, oscillations for  $\delta \neq 0$ . The ability to vary the detuning, phase and amplitude of the resonant laser pulses enables multi-axis control of the optical qubit [37, 56].

We next leverage multi-axis control to probe the coherence of the optical transition directly through pulsed

resonant excitation. Figure 3(a) displays the measurements of Ramsey interferometry and Hahn-echo dynamical decoupling. We read out the state of the emitter by integrating the fluorescence after the final  $\pi/2$ -rotation. The population contrast between the ground and excited states is measured by varying the phase of this  $\pi/2$  pulse [33, 56]. The solid teal curve in Fig. 3(a) is a fit to the Ramsey-interferometry data using the master-equation model employed previously [37]. This yields an inhomogeneous dephasing time  $T_2^*$  of  $4.54(2)$  ns which corresponds to a shot-to-shot spectral drift of  $82.6(5)$  MHz at full width at half maximum (FWHM). The Hahn-echo contrast decay envelope provides a measurement of the pure-dephasing rate  $\Gamma_{PD}$ . A fit to the raw contrast results in  $\Gamma_{PD}/2\pi = 6.39(14)$  MHz and an inferred homogeneous linewidth of  $\Gamma/2\pi = (\Gamma_0 + 2\Gamma_{PD})/2\pi = 34.8(7)$  MHz. This linewidth is only a factor of  $1.63(4)$  from its intrinsic Fourier limit ( $\Gamma_0$ ) despite the emitter's close proximity ( $< 60$  nm) to nanostructured surfaces.

We next probe the spectral stability on longer timescales by observing the shot-to-shot variation of the transition frequency through fast photoluminescence excitation (PLE) scans at  $s = 0.8(1)$ . Each scan alternates between 500 ms of resonant 619-nm and off-resonant 532-nm excitation. During the resonant section,  $\delta/2\pi$  is repeatedly scanned from  $-100$  MHz to  $100$  MHz using  $3$ - $\mu$ s-long linearly-chirped laser pulses. Figure 3(b) shows the spectral evolution over 34 minutes. Within each vertical cut, and after correcting for power broadening, the linewidth is  $35(10)$  MHz, commensurate with the homogeneous linewidth inferred from the Hahn-echo measurement. This indicates that there are no significant, additional, dephasing mechanisms present between the tens of nanoseconds timescale probed in Fig. 3(a), and the 500-ms timescale probed by the PLE scans in Fig. 3(b).

The central frequency does not show significant variation over consecutive scans. Over the entire measurement it tends towards a normal distribution with a FWHM of  $64(1)$  MHz. Such inhomogeneous broadening has been observed for the nitrogen-vacancy center [14], and more recently the SnV center [35], and is likely due to the off-resonant laser rearranging the local charge environment. For the nitrogen-vacancy center, the photo-induced spectral broadening can be over  $0.5$  GHz, or 35-times broader than the intrinsic linewidth [59]. In contrast, electric noise only detunes the SnV center through a second-order Stark shift [25] leading to a significantly narrower inhomogeneous linewidth. The deviation with respect to the Ramsey measurement could be attributed to the variation in 532-nm irradiation time. Figure 3(b) shows an inhomogeneous broadening of a factor of  $\sim 2$  compared to its homogeneous linewidth.

The 532-nm off-resonant excitation photo-stabilizes the emitter as resonant excitation pumps the emitter stochastically into a dark state [35]. To establish the timescale of this pumping, we monitor the decay of PLE intensity following a 532-nm off-resonant pulse. Repeatedly scanning  $300$  MHz across the transition ensures in-

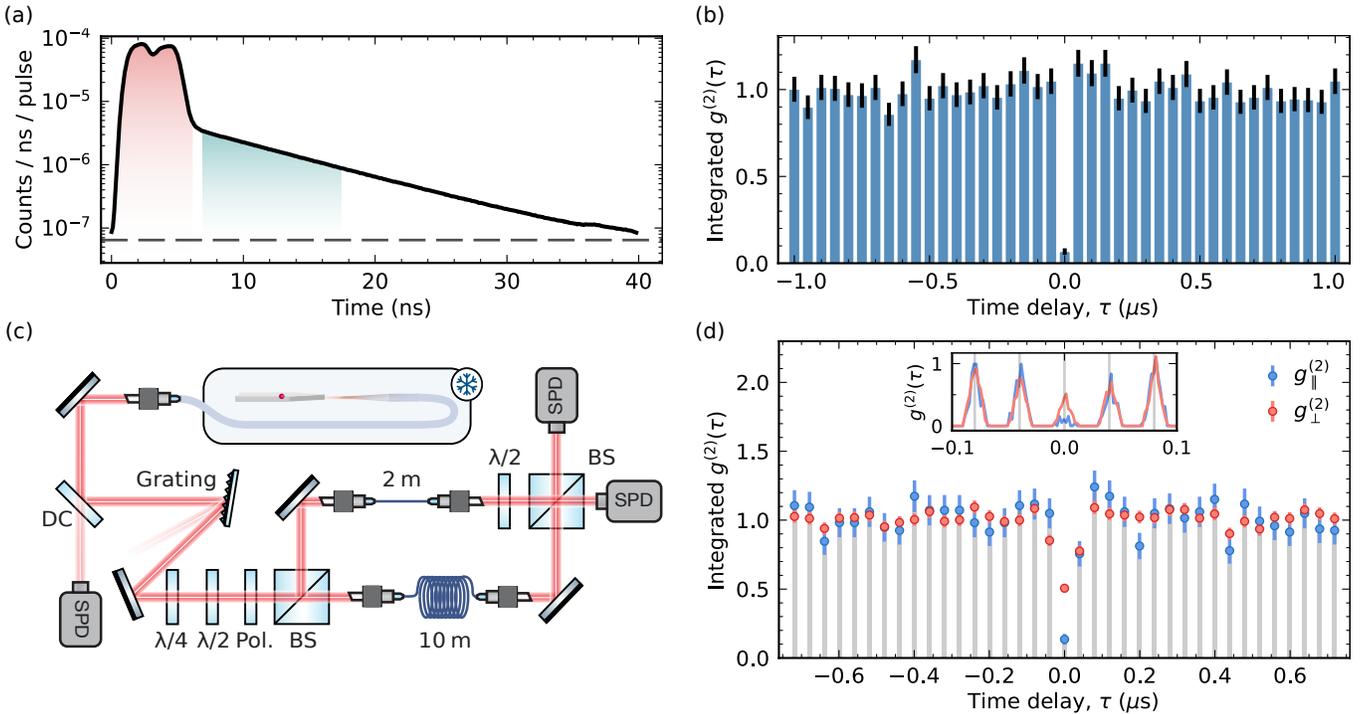


FIG. 4. (a) Time-resolved ZPL fluorescence under excitation with a 4-ns-long resonant laser pulse, at  $s = 45(20)$ . Red and teal shadings correspond to applied laser excitation and ZPL fluorescence collection, respectively. The average background (dashed line) is extracted with a far-detuned ( $\delta/2\pi = 6$  GHz) laser pulse [37]. (b) Second-order autocorrelation of the ZPL fluorescence under pulsed resonant excitation. The recorded value  $g^{(2)}(0)_{\text{raw}} = 0.067(23)$  corresponds to a background-corrected photon purity of  $99.7^{+0.3}_{-2.5}\%$ . (c) Experimental setup for the HOM measurement. Abbreviations: quarter-wave plate ( $\lambda/4$ ), half-wave plate ( $\lambda/2$ ), polarizer (Pol.), beamsplitter (BS) and dichroic mirror (DC). (d) Pulsed two-photon interference measurement for the parallel (blue) and perpendicular (red) polarization configurations, where the coincidence counts are integrated for each pulse. The background-corrected photon indistinguishability is  $73(13)\%$  for photons collected within the teal region of panel a. Inset: time-resolved distribution of coincidences around  $\tau = 0$ .

sensitivity to spectral diffusion. Figure 3(c) shows the decay of PSB fluorescence for various resonant excitation powers. The solid curves in Fig. 3(c) are fits to a bi-exponential model, where the dominant component is the fast timescale of pumping into the dark state. The slow timescale contributes 10(1)% of the total fluorescence at the highest power, likely due to a residual photo-active background [60, 61]. The dark state pumping rate increases linearly with  $s$ , in accordance with previous reports [35]. At the highest laser power, fluorescence persists over milliseconds allowing  $\sim 10^6$   $\pi$ -pulses to be applied before pumping into the dark state. When considered in concert with Fig. 3(b), our emitter remains spectrally stable and optically active over many milliseconds, as required for quantum networking [62].

To probe the indistinguishability of the emitted photons, we isolate the ZPL using spectral filtering. Residual resonant laser is suppressed by  $>30$  dB by pulsing the resonant excitation with two electro-optic modulators and time-tagging ZPL fluorescence after resonant excitation. Figure 4(a) shows the time-resolved ZPL fluorescence after a 4-ns resonant excitation pulse. The 90 dB total laser suppression yields a signal-to-background ratio of  $23.91(3)$  for the ZPL fluorescence collected within 11.1

ns after the excitation pulse [37].

We measure a zero-delay second-order intensity correlation  $g^{(2)}(0)_{\text{raw}}$  of  $0.067(23)$ , as shown in Fig. 4(b) [63]. This is comparable to previous reports where only the PSB was collected, despite the added technical challenge of resonant collection in this work [33, 34]. After correcting for detector dark counts, the background-corrected photon purity  $(1 - g^{(2)}(0))$  is  $99.7^{+0.3}_{-2.5}\%$ .

Figure 4(c) shows our HOM-interferometry setup [37]. Before the interferometer, a dichroic mirror and a grating filter photons within a 245(6) GHz bandwidth (FWHM) centered on the ZPL transition. This filtered emission is routed through polarization control optics and into the interferometer with a relative time delay  $\Delta t$  of 39.93(2) ns between the two arms. To ensure temporal overlap between subsequently emitted photons we apply laser pulses with a repetition period matching  $\Delta t$ . A half-wave plate placed in the short arm of the interferometer controls the relative polarization between the two interfering photons.

Figure 4(d) shows the second-order intensity correlation measured across the two output ports of the interferometer for ZPL fluorescence integrated over the first 11.1 ns after the excitation pulse. When the polarizations

of two interfering photons are orthogonal to each other, we measure  $g_{\perp}^{(2)}(0) = 0.51(3)$ , in agreement with the theoretically expected value of 0.5. When the polarizations of the photons are matched, we measure  $g_{\parallel}^{(2)}(0) = 0.22(3)$ . The raw visibility,  $V_{\text{raw}} = 1 - g_{\parallel}^{(2)}(0)/g_{\perp}^{(2)}(0) = 56(8)\%$ , is a measurement of the indistinguishability of the emitted photons. Correcting for the background and finite classical interferometric contrast, we calculate a photon indistinguishability  $V$  of 73(13)%. When the time window is extended to the full decay observed in Fig 4(a) the indistinguishability is reduced to  $V = 63(9)\%$  [37].

We compare our measured indistinguishability to a model that includes homogeneous broadening as the source of photon distinguishability. Given the measured  $\Gamma_{\text{PD}}$ , the expected HOM visibility  $V_{\text{sim}}$  is 63.2(4)% in agreement with our measured value [37]. Considering the slow spectral diffusion measured in Fig 3(b), the HOM visibility measured here should extend to interferometer delays greater than 100 ms.

In this work, we demonstrate coherent control of the optical transition of the SnV in diamond. We achieve a  $\pi$ -pulse time of 1.71(1) ns and two-photon quantum interference of resonant photons with an indistinguishability of 63(9)%. Purcell enhancement of the emitter, such as through optical microcavities [64] or photonic crystal cavities [28–30], increases the radiative decay rate and thus reduces the sensitivity to dephasing. An overall Purcell enhancement factor of 12, already achieved using SnV centers in nanocavities [29, 30], would yield a photon indistinguishability in excess of 95%. Moreover, fabrication and material processing improvements, known to result in improved optical quality of SnV centers in bulk diamond [22], should enhance the optical coherence. These advancements provide a feasible route to near-unity photon indistinguishability in the near-term.

Using the two-photon interference presented in this work, in combination with control of the spin degree of freedom [33], one could realize spin-photon entanglement [5, 65, 66] and entanglement of remote emitters [4, 67, 68]. Modest improvements in the indistinguishability of the photons would position multi-photon entangled states as the next achievement using the SnV center [69, 70].

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

\* These authors contributed equally to this work.

† Correspondence should be addressed to: dorian.gangloff@eng.ox.ac.uk, englund@mit.edu, ma424@cam.ac.uk.

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