Experimental observation of optical Weyl points and Fermi arc-like surface states

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Weyl fermions are hypothetical two-component massless relativistic particles in three-dimensional (3D) space, proposed by Hermann Weyl in 1929. Their band-crossing points, called ‘Weyl points’, carry a topological charge and are therefore highly robust. There has been much excitement over recent observations of Weyl points in microwave photonic crystals and the semimetal TaAs. Here, we report on the experimental observation of ‘type-II’ Weyl points of light at optical frequencies, with the photons having a strictly positive group velocity along one spatial direction. We use a 3D structure consisting of laser-written waveguides, and show the presence of type-II Weyl points by observing conical diffraction along one axis when the frequency is tuned to the Weyl point; and observing the associated Fermi arc-like surface states. The realization of Weyl points at optical frequencies allows these novel electromagnetic modes to be further explored in the context of linear, nonlinear, and quantum optics.

The observation of Weyl points, in microwave photonics1 and condensed matter physics3, has attracted a great deal of attention because they constitute the simplest possible topologically nontrivial band structure in three dimensions. Their topological protection guarantees the existence of ‘Fermi arc’ surface states4 in solid-state materials, and they can be associated with many interesting phenomena, including chiral anomalies5, unconventional superconductivity6, and large-volume single-mode lasing7. They take two distinct forms: type-I Weyl points have a Fermi surface (in the electronic context) or isofrequency surface (in the photonics context) that is point-like, whereas type-II Weyl points have a Fermi or isofrequency surface that is conical8–11. In photonics, type-I Weyl points were predicted12,13 and subsequently observed14 in macroscopic photonic crystals at microwave frequencies. There have been significant efforts, both theoretical14–20 and experimental21–23, to realize Weyl points at the technologically important optical frequency regime; this is challenging, however, due to the need for 3D fabrication of intricate photonic crystal structures. Photonic type-II Weyl points have also been previously theoretically proposed24,25, although not experimentally observed.

Here we observe photonic type-II Weyl points, at optical frequencies, in a 3D photonic crystal structure of evanescently coupled single-mode waveguides, fabricated using femtosecond direct laser writing26 (see Methods). The waveguides form an array aligned along a particular spatial axis z, as shown in Fig. 1a. For an appropriately designed array, we show that the 3D photonic band structure is locally described by a \(2 \times 2\) Weyl Hamiltonian

\[
\hat{H} \approx v_x (\hat{\sigma}_x + \Delta z \hat{\sigma}_z) + v_z \Delta z \hat{\sigma}_z (\hat{I} - |b| \hat{\sigma}_z)
\] (1)

whose eigenvalues \(\Delta \omega\) are the band frequencies relative to a chosen origin (the Weyl frequency). Here \(\Delta z = (\Delta k_x, \Delta k_z)\) and \(\Delta k_z\) are the wavevector components transverse and parallel to the waveguide axis respectively (relative to the Weyl point), with group velocities \(v_x, v_z\). \(\hat{\sigma}_x, \hat{\sigma}_y, \hat{\sigma}_z\) are the Pauli matrices, \(\hat{I}\) is the identity operator, and \(b\) is a dimensionless parameter. For our system, \(|b| \ll 1\). Equation (1) is derived by taking the standard ‘paraxial’ description of weakly confined optical waveguide modes26,27, and casting the results into a 3D photonic band structure; details are provided in the Supplementary Information.

The Hamiltonian (1) describes a type-II Weyl point13,24. As shown schematically in Fig. 1b, the dispersion is strongly anisotropic, with both bands having positive group velocities in the \(z\) direction. This is because the waveguide modes move in a single direction along \(z\) with negligible backscattering, even as they undergo diffraction in the transverse directions \(x\) and \(y\). The corresponding band structure has distinctive isofrequency surfaces: as shown in Fig. 1b, the isofrequency surfaces at \(\Delta \omega = 0\) are cones in \(k\) space, whereas for small nonzero \(\Delta \omega\) the isofrequency surfaces become hyperboloids. We will use this fact in our experimental probe of the optical Weyl point.

Weyl points are topologically protected because the Pauli matrices, together with \(\hat{I}\), span the space of \(2 \times 2\) Hermitian matrices. As long as coupling to any other modes of the system can be neglected, within this two-mode subspace the Weyl point acts as a source or sink of Berry curvature, carrying a quantized topological charge \(C = \pm 1\). Under weak perturbations, a Weyl point merely moves in \(\Delta \omega\) and \(k\) space; it can be eliminated only via annihilation with a partner Weyl point of opposite charge. The Hamiltonian (1) with charge \(C = -1\) can be modified to include a partner by replacing the \(|b| \hat{\sigma}_z\) term with \(|b|(1 + \Delta k_z / q) \hat{\sigma}_z\); the partner Weyl point with opposite charge \(C = +1\) appears at \(\Delta k_z = -q\) and at a lower frequency, as shown in Fig. 1c. At spatial boundaries, each such pair of Weyl points is associated with Fermi arc-like surface states, whose isofrequency dispersion curves form open arcs connecting the projections of the two Weyl points in the surface Brillouin zone25. (These states are the bosonic counterparts of the Fermi arc states observed in solid-state materials with Weyl points.) Importantly, the surface states span the.

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range of frequencies separating the type-II Weyl points, as shown in Fig. 1d. This will provide us with another experimental probe for the optical Weyl point.

Our waveguide array design, depicted in Fig. 1a, consists of a bipartite square lattice with two helical waveguides in each unit cell, both having clockwise helicity, radius $R$, lattice constant $a$, and period $Z$ in the $z$ direction. A microscope image of a cross-sectional cut of the array at the input facet is shown in Fig. 2a. Adjacent waveguides are out of phase by a half-cycle, so that they do not evolve in synchrony, but their nearest-neighbour distance changes through a period in $z$ (ref. 28); the design is closely linked to the concept of anomalous Floquet topological insulators\textsuperscript{30}. The helical modulation breaks the inversion symmetry of the structure—a necessary condition for realizing Weyl points in time-reversal-invariant systems such as this one\textsuperscript{14}. The time-reversal symmetry $T$ implies the Brillouin zone must contain at least four Weyl points: for each Weyl point, $T$ maps to an equivalent Weyl point of the same charge located at the opposite side of the Brillouin zone, corresponding to modes propagating in the opposite direction along $z$.

As mentioned above, waveguide arrays are commonly analysed in the paraxial limit\textsuperscript{36}, by splitting the three spatial dimensions into the transverse $(x, y)$ plane and the axial direction $z$. The Maxwell equations describing the diffraction of light through the array can be reduced to a two-dimensional (2D) Schrödinger-like equation, with $z$ acting as a temporal direction. The operating wavelength (or frequency) enters as an adjustable parameter in the equation, while the wavevector component in the $z$ direction, $k_z$, acts as the effective ‘energy’. In other words, the paraxial equation describes isofrequency surfaces of the full 3D photonic band structure. Crucially, surfaces taken at different frequencies may correspond to topologically distinct 2D band structures of the paraxial equation. In the Supplementary Information, we show that the isofrequency surfaces for $\delta \omega > 0$ (or $\delta \omega < 0$) correspond to conventionally insulating (topologically insulating) 2D paraxial band structures; the isofrequency surface at the Weyl point, $\delta \omega = 0$, corresponds to a 2D paraxial band structure poised at a topological transition. Notably, the $\delta \omega < 0$ regime is also where the Fermi arc-like surface states emerge.

Previously, a honeycomb lattice of helical waveguides was used to realize a photonic topological insulator in ref. 30; however, that design does not readily support observable optical Weyl points. The reason why the present design is ideal for supporting Weyl
Figure 2 | Theoretical and numerical demonstration of topological phase transition associated with type-II Weyl points. a. Microscope image of the output facet of structure, representing a two-dimensional cut of the waveguide array for fixed $z$. b. Numerically determined phase diagram of the structure, as a function of lattice constant $a$ and wavelength $\lambda$. Type-II Weyl points reside along the red curves, and Fermi arc-like surface states exist between these two curves (yellow region). c. Bulk band structure for the two relevant bands plotted as a function of $k_z$ (in the $k = k_y = 0$ plane, using the extended-zone scheme). Type-II Weyl points arise at their intersection. d-f. Isofrequency surfaces for the topologically trivial case (no Fermi arc-like states), at the Weyl point (WP), and the topological (with Fermi arc-like states) case, at $a = 29, 27$ and $25 \mu m$, at wavelengths $1,450$ nm, $1,525$ nm and $1,600$ nm, respectively. The open circles in the phase diagram shown in b correspond to the band structures in d-f. All results in b-f are calculated numerically\textsuperscript{28}, using experimental parameters.

points can be described as follows. The helical waveguides in the two sublattices come closer and farther apart from one another over the course of propagation along $z$. Therefore, the effect of waveguide helicity is strongest when either the lattice constant is small (so the helices get very close together and the modes are strongly coupled), or the wavelength is sufficiently long (such that the individual waveguide modes are large and therefore are strongly coupled). Therefore, the lattice is trivial if the waveguides are weakly coupled, and topological if they are strongly coupled. The transition takes place in the middle—and can be tuned to be straightforwardly accessed experimentally in the paraxial regime. The Weyl point occurs at the transition.

We analyse the presence of the Weyl point in systems of different lattice constant, $a$. Figure 2b shows a phase diagram versus $a$ and wavelength $\lambda$ which indicates the lines in this parameter space for which Weyl points appear. The lines indicate the degeneracies in the bulk band structure (for example, the 1D band structure cut in Fig. 2c, plotted in the extended-zone scheme for clarity) dividing the topological regime (which supports Fermi arc-like states) from the trivial regime (which does not). Note that while changes in a shift the Weyl points in frequency and wavevector, the Weyl points exist for all $a$ in this range. Their robustness under this smooth deformation is a signature of their topological protection. In Fig. 2d-f, we plot three isofrequency surfaces calculated for structures of radius $R = 4 \mu m$ and period $Z = 1$ cm: panel d shows the $\delta \omega > 0$ case (lattice constant $a = 29 \mu m$, wavelength $\lambda = 1,450$ nm), panel e shows the $\delta \omega = 0$ case ($a = 27 \mu m$, $\lambda = 1,525$ nm); and panel f shows the $\delta \omega < 0$ case ($a = 25 \mu m$, $\lambda = 1,600$ nm). These numerical results were obtained via the evolve-and-project method described in ref. 28. These calculations show that the Weyl point resides at $k = (0, 0, 0.08\pi/Z)$ at $\lambda = 1,525$ nm for $a = 27 \mu m$, where $k$ is the wavevector in the three-dimensional Brillouin zone of the photonic crystal band structure.

We prove the existence of a type-II Weyl point in two distinct experiments: by observing conical diffraction associated with the isofrequency surface of a type-II Weyl point\textsuperscript{13}, and by observing the Fermi arc-like surface states that emerge from the Weyl point.

The conical diffraction experiment relies on the fact that the isofrequency surface for the type-II Weyl point is a cone, as shown in Figs 1b and 2e. At this frequency, every Bloch wave has the same group velocity $v_z$ in the transverse plane. Thus, an initial wavepacket injected into the waveguide array, which can be decomposed into a superposition of Bloch waves, evolves into a ring as depicted in Fig. 3a; the angle at which the ring expands is given by $v_z / v_g$, where $v_g$ is the group velocity in the $z$ direction. In the context of paraxial optics, this is precisely conical diffraction, marked by a fixed angle of diffraction relative to the axis of the waveguide array\textsuperscript{31–32}.

We emphasize that a type-I Weyl point would not exhibit conical diffraction, because the associated isofrequency surface is a point rather than a cone. Moreover, for a given photonic band structure with a type-II Weyl point, conical diffraction occurs at a single probe wavelength; for $\delta \omega \neq 0$, the isofrequency surfaces are hyperboloids rather than cones, and hence an injected wavepacket produces a broad diffraction pattern rather than a ring. This wavelength selectivity is notably distinct from the conical diffraction phenomena observed for band structures lacking type-II Weyl points, such
Figure 3 | Conical diffraction as a signature of the existence of type-II Weyl points. a, Schematic diagram of conical diffraction occurring in the waveguide array at a type-II Weyl point. b–d, Intensity plots at the output facet, as we sweep through the Weyl point, at \( a = 29, 27 \) and \( 25 \) \( \mu m \), and wavelengths \( 1,450 \text{ nm}, 1,525 \text{ nm} \) and \( 1,600 \text{ nm} \), respectively. The green circles indicate the position of the input waveguides. Clear conical diffraction is observed in c, at the Weyl point. e–g, Full-wave simulations corresponding to the parameters of b–d. h–j, Plot of the quantity \( C \), obtained from experimental data, as a function of \( \lambda \), which quantifies how ring-like the wavefunction for \( a = 29, 27, 25 \) \( \mu m \), as a function of wavelength from \( 1,450–1,650 \) \text{ nm}. In h and j, there is no clear minimum, indicating the lack of a Weyl point within this wavelength range. However, i shows a clear minimum, at the wavelength where the Weyl point lies. This minimum corresponds to the wavefunction shown in c and f.

We perform the conical diffraction experiment by injecting light into a single waveguide at the input facet, which then couples through a waveguide splitter to a pair of neighbouring waveguides (within one unit cell at the centre of the lattice) with equal intensity and phase. The in-coupling region occupies the first 1 cm of the chip. The resulting diffraction patterns, imaged at the end of the chip, are shown in Fig. 3b–d. For lattice constants \( a = 29 \) \( \mu m \) (at wavelength 1,450 nm, Fig. 3b) and \( a = 25 \) \( \mu m \) (at wavelength 1,600 nm, Fig. 3d), for which the isofrequency band structure is in a conventional and topological phase, respectively, we observe a filled-in disc-like
diffraction pattern, which is characteristic of parabolic dispersion. When the lattice constant is tuned to $a = 27 \mu m$, we observe a clear ring-like conical diffraction pattern at $\lambda = 1,525 \text{ nm}$, shown in Fig. 3c. Deviations from a perfect ring occur because the lattice is discrete, which makes it impossible for the light to reside on a perfectly circular and zero-width ring. Furthermore, the diffraction pattern is distorted by the square symmetry of the lattice, which induces a saddle point in the isofrequency surface at the Brillouin zone edge. Figure 3e–g shows full-wave beam-propagation simulations, which agree strongly with the experimental results of Fig. 3b–d. Thus, the presence of conical diffraction clearly establishes the presence of the Weyl point.

To quantify the observation of conical diffraction, we define a dimensionless parameter that measures the degree to which the diffraction pattern is conical:

$$C = \frac{\int dr r^2 |\psi(r)|^2}{(\int dr |\psi(r)|^2)^2}$$

where $r = (x,y)$ and $r = \sqrt{x^2 + y^2}$ is the distance from the origin (the origin is defined to be the centre point between the two excited waveguides). The quantity $C$ measures how ‘ring-like’ a wavefunction is, and is reminiscent of the inertial moment of a rotating body in a mechanical system (but is dimensionless in this case). It takes the value 1 for an infinitely thin ring, and is larger than 1 for all other patterns. In Fig. 3h–j we plot the experimentally obtained $C$ against wavelength for the three previous lattice constants, namely $a = 29, 27, 25 \mu m$. The monotonic behaviour of Fig. 3h and j shows that for lattice constants $a = 29\mu m$ and $25\mu m$ the Weyl point does not occur within the wavelength range of the tunable laser (1,450–1,650 nm), or is very near the boundary. That said, the Weyl points do exist within the spectrum for each value of $a$ described here. The minimum in $C$ observed in Fig. 3i corresponds to conical diffraction (see Fig. 3c and f for experimental and numerically computed conical diffraction patterns). The conical diffraction associated with the isofrequency surface provides direct evidence of the existence of

**Figure 4 | Direct observation of Fermi arc-like surface states.** a. Schematic diagram of light confinement to the surface as a result of a Fermi arc-like surface state. b. Position of the four type-II Weyl points in the 3D cubic Brillouin zone (represented by the enclosing box), with their topological charges indicated. WP1 and WP2 are partner Weyl points, and WP3 and WP4 are their time-reversed equivalents (exact values of $k$ in text). c. Plot of the Fermi arc-like dispersion relation in the surface Brillouin zone, within the range of experimentally accessible wavelengths, as a function of $k_x$ and $k_z$ (the surface is terminated in the $y$-direction). This is calculated using the method of ref. 28. d–h. Output intensity plots, when light is input at the centre of the top surface of the structure (indicated by green circles) at wavelength 1,550 nm, for decreasing $a = 29, 28, 27, 26, 25 \mu m$. For decreasing lattice constant $a$, increased confinement to the surface indicates the formation and presence of Fermi arc-like surface states. i–m. Corresponding full-wave beam propagations, showing strong agreement with experimental results. n–r. Numerically calculated isofrequency contours, showing the presence of surface states forming at $a = 27 \mu m$ as the Weyl point is crossed. The red and blue curves indicate surface states on the top and bottom of the sample, respectively. The trajectories of the surface state wavepackets are indicated by red arrows.
the type-II Weyl point. As mentioned above, for every Weyl point there must be a partner Weyl point—which in this case occurs at \( k = (0, 0, 0.53\pi/Z) \) at wavelength \( \lambda = 2,113 \text{ nm} \) for \( a = 27 \text{ nm} \), and lies outside the accessible parameter range of the experiment.

Next, we demonstrate the existence of the Fermi arc-like surface states associated with the type-II Weyl point. As stated earlier, these surface states span the range of frequencies between this Weyl point and its partner. Specifically, we expect them to appear below the Weyl frequency (where the 2D paraxial band structure becomes topologically nontrivial) and link the two bulk bands. To probe for surface states, we inject light via a single waveguide at the top of the lattice and observe the output facet. If a surface state is present, light should stay confined to the surface (see schematic depiction in Fig. 4a); otherwise it would diffract into the bulk. The positions of all Weyl points and corresponding Fermi arc-like states in the Brillouin zone are shown in Fig. 4b—note that we have identified a total of four, but focus on the single Weyl point associated with the type-II Weyl point. As stated earlier, these surface states span the range of frequencies between this Weyl point and its partner.

In conclusion, we have made direct experimental observations of Weyl fermion semimetals and topological Fermi arcs in the electronic structure of pyrochlore iridates. Phys. Rev. B 83, 205101 (2011).

Xu, G., Weng, H., Wang, Z., Dai, X. & Fang, Z. Chern semimetal and the quantized anomalous Hall effect in HfGe2Se4, Phys. Rev. Lett. 107, 186806 (2011).

Burkov, A. A. & Balents, L. Weyl semimetal in a topological insulator multilayer. Phys. Rev. Lett. 107, 127205 (2011).

Potter, A. C., Kimchi, I. & Vishwanath, A. Quantum oscillations from surface Fermi arcs in Weyl and Dirac semimetals. Nat. Commun. 5, 5161 (2014).

Nieslen, H. B. & Ninomiya, M. The Adler–Bell–Jackiw anomaly and Weyl fermions in a crystal. Phys. Lett. B 130, 389–396 (1983).

Burkov, A. A. Chiral anomaly and transport in Weyl metals. J. Phys. Condens. Matter 27, 113201 (2015).

Cho, G. Y., Bardarson, J. H., Lu, Y. M. & Moore, J. E. Superconductivity of doped Weyl semimetals: finite-momentum pairing and electronic analog of the \( ^3 \text{He} \text{-A} \) phase. Phys. Rev. B 86, 214514 (2012).

Bravo-Abad, J., Joannopoulos, J. D. & Soljacic, M. Enabling single-mode behavior over large areas with photonic Dirac cones. Proc. Natl Acad. Sci. USA 109, 9761–9765 (2012).

Xu, Y., Zhang, C. & Zhang, C. Structured Weyl points in spin-orbit coupled fermionic superfluids. Phys. Rev. Lett. 115, 263504 (2015).

Soluyanov, A. A. et al. Type-II Weyl semimetals. Nature 527, 495–498 (2015).

Lu, L., Fu, L., Joannopoulos, J. D. & Soljacic, M. Weyl points and line nodes in gyroid photonic crystals. Nat. Photon. 7, 294–299 (2013).

Wang, L., Jian, S.-K. & Yao, H. Topological photonic crystal with equifrequency Weyl points. Phys. Rev. A 93, 061801 (2016).

Bravo-Abad, J., Lu, L., Fu, L., Boljan, H. D. & Soljacic, M. Weyl points in photonic-crystal superlattices. 2D Mater. 2, 034013 (2015).

Gao, W. et al. Photonic Weyl degeneracies in magnetized plasma. Nat. Commun. 7, 12435 (2016).

Xiao, M., Chen, W.-J., He, W.-Y. & Chan, C. T. Synthetic gauge flux and Weyl points in acoustic systems. Nat. Phys. 11, 920–924 (2015).

Yang, Z. & Zhang, B. Acoustic Weyl nodes from stacking dimerized chains. Phys. Rev. Lett. 117, 224301 (2016).

Xiao, M., Lin, Q. & Fan, S. Hyperbolic Weyl point in reciprocal chiral metamaterials. Phys. Rev. Lett. 117, 057401 (2016).

Chen, W.-J., Xiao, M. & Chan, C. Photonic crystals possessing multiple Weyl points and the experimental observation of robust surface states. Nat. Commun. 7, 13038 (2016).

Peng, S. et al. Three-dimensional single gyroid photonic crystals with a mid-infrared bandgap. ACS Photon. 3, 1131–1137 (2016).

Peng, S. et al. Gyroid photonic crystal with Weyl points: synthesis and mid-infrared photonic characterization. APS March Meeting 2016, Abstract No52.013 (2016).

Xiao, M., Lin, Q. & Fan, S. Hyperbolic Weyl point in reciprocal chiral metamaterials. Phys. Rev. Lett. 117, 057401 (2016).

Szameit, A. & Noile, S. Discrete optics in femtosecond-laser-written photonic structures. J. Phys. B 43, 163001 (2010).

Yariv, A. & Yeh, P. Optical Waves in Crystals Vol. 10 (Wiley, 1984).

Fleischer, J. W., Segov, M., Efremidis, N. K. & Christodoulides, D. N. Observation of two-dimensional discrete solitons in optically induced nonlinear photonic lattices. Nature 422, 147–150 (2003).

Levký, D., Rechtsman, M. C. & Chong, Y. D. Anomalous topological phases and unspaired Dirac cones in photonic Floquet topological insulators. Phys. Rev. Lett. 117, 013902 (2016).

Rudner, M. S., Lindner, N. H., Berg, E. & Levin, M. Anomalous edge states and the bulk-edge correspondence for periodically driven two-dimensional systems. Phys. Rev. X 3, 031005 (2013).

Rechtsman, M. C. et al. Photonic Floquet topological insulators. Nature 496, 209–210 (2013).

Berry, M., Jeffreys, M. & Lanney, J. Conical diffraction: observations and theory. Proc. Roy. Soc. A 662, 1629–1642 (2006).

Peleg, O. et al. Conical diffraction and gap solitons in honeycomb photonic lattices. Phys. Rev. Lett. 98, 103901 (2007).

Kim, Y., Wieder, B. J., Kane, C. L. & Rappe, A. M. Dirac line nodes in inversion-symmetric crystals. Phys. Rev. Lett. 115, 036806 (2015).

Levký, D. & Chong, Y. D. Edge solitons in nonlinear–photonic topological insulators. Phys. Rev. Lett. 117, 143901 (2016).

Zhen, B. et al. Spawning rings of exceptional points out of Dirac cones. Nature 525, 354–358 (2015).

Peruzzo, A. et al. Quantum walks of correlated photons. Science 329, 1500–1503 (2010).

Rechtsman, M. C. et al. Topological protection of photonic path entanglement. Optica 3, 925–930 (2016).
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Author contributions
J.N. carried out experimental measurements and performed the data analysis; S.H. developed the laser fabrication process and characterized the samples under the supervision of K.P.C. and with guidance from M.C.R.; D.L., C.Y.D. and M.C.R. conceived the idea and performed theoretical analysis and calculations; M.C.R. supervised the project.

Additional information
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Competing financial interests
The authors declare no competing financial interests.
Methods

The waveguides are written in Corning Eagle XG borosilicate glass, refractive index $n_0 = 1.473$. We employed a titanium:sapphire laser and amplifier system (Coherent: RegA 9000) with pulse duration 270 fs, repetition rate 250 kHz, and pulse energy 950 nJ. The laser writing beam was sent through a beam-shaping cylindrical telescope to control the shape and size of the focal volume. The beam was then focused inside the glass chip using a $\times 80$, aberration-corrected microscope objective (NA = 0.75). A high-precision three-axis Aerotech motion stage (model ARL20020) is used to translate the sample during fabrication. Experiments are performed by butt-coupling a single-mode optical fibre to waveguides at the input facet of the chip, which subsequently couples to the waveguide array. The input light is supplied by a tunable mid-infrared diode laser (Agilent 8164B), which can be tuned through the 1,450–1,650 nm wavelength range. After a total propagation distance of 4 cm within the array, the light output from the waveguide array is observed using a 0.2 NA microscope objective lens and a near-infrared InGaAs camera (ICI systems).

Data availability. The data that support the plots within this paper and other findings of this study are available from the corresponding author upon reasonable request.