

Supplementary information for
Mode Vernier photoacoustic spectroscopy gas sensing using 3D-
microprinted on hollow-core fiber tip

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Note 1. Theoretical Derivation of Effective Refractive Index Difference Between LP₀₁ and LP₁₁ Modes in BHCF

This supplementary material presents a rigorous theoretical derivation of the effective refractive index (RI) difference between the LP₀₁ fundamental mode and LP₁₁ higher-order mode in the BHCF investigated in this work, based on the classic asymptotic matrix theory and leaky mode model for Bragg fibers¹⁻². The derivation fully accounts for the structural parameters of the fabricated BHCF, and validates the physical origin of the experimentally observed giant effective RI difference of 4.8×10^{-2} between the two modes.

The BHCF features a 32 μm air hollow core ($n_{co} = 1.0$ at $\lambda_0 = 1550$ nm), a fused silica background cladding ($n_L = 1.444$ at 1550 nm), and four embedded high-RI annular Bragg layers ($n_H = 1.454$ at 1550 nm)³. Each high-RI layer has a radial thickness of $l_H = 1.06$ μm , with a 3.07 μm -thick low-RI fused silica gap (l_L) between adjacent high-RI layers, giving a single Bragg period of $L_a = l_H + l_L = 4.13$ μm , as shown in Fig. S1a. For the theoretical derivation, the following fundamental constants and parameters are adopted: vacuum wave vector $k_0 = 2\pi/\lambda_0 = 4.0536$ μm^{-1} , low-index dielectric constant $\varepsilon_L = n_L^2 = 2.0851$, and high-index dielectric constant $\varepsilon_H = n_H^2 = 2.1141$.

Numerical simulations based on the finite element method (FEM) confirm that the BHCF supports two low-order eigenmodes guided within its air hollow core, namely the fundamental LP₀₁ mode and the first higher-order LP₁₁ mode, as defined in standard cylindrical waveguide optics¹⁻². The simulated mode field distributions of the LP₀₁ and LP₁₁ modes at the 1550 nm operating wavelength are presented in Fig. S1b and S1c, respectively. Notably, the LP₁₁ mode acts as a strongly leaky mode in the 4-period weak-modulation Bragg grating of the BHCF. At the same time, it still maintains a well-defined, stable double-lobe mode field profile and a deterministic effective RI within the millimeter-scale device length employed in this work. The effective RI difference between the two modes is numerically calculated to be $\sim 4.8 \times 10^{-2}$, which is consistent with the experimental measurement. The beat length l_b between the two co-propagating modes, defined as the propagation distance for the two modes to accumulate a 2π phase difference, is given by:

$$l_b = \frac{\lambda_0}{\Delta n_{eff}} \quad \text{S(1)}$$

where $\lambda_0 = 1550$ nm is the operating wavelength, and $\Delta n_{eff} = 4.8 \times 10^{-2}$ is the effective RI difference between the two modes. Substituting the parameters into Eq. (S1) yields a beat length of ~ 32 μm , which matches the experimental characterization.

All derivations in this work are based on the fundamental solutions of Maxwell's equations for cylindrical symmetric waveguides, with core definitions and theoretical

frameworks detailed below. For a guided mode propagating along the fiber axis (z-axis), the effective refractive index is strictly defined as:

$$n_{eff} = \frac{\beta_r}{k_0} \quad S(2)$$

where β_r is the real part of the mode propagation constant, which determines the phase propagation characteristics of the mode. Unlike conventional step-index fibers that guide light via total internal reflection, BHCFs confine light through periodic Bragg reflection, meaning all guided modes are inherently leaky modes with a complex propagation constant¹:

$$\beta = \beta_r + i\beta_i \quad S(3)$$

The imaginary part β_i governs the leakage loss of the mode, where the optical power follows the attenuation $P(z)=P_0e^{-2\beta_i z}$. The conversion relationship between the loss in decibel (dB) units and β_i is:

$$\alpha_{dB/m} = 8.6858\beta_i \quad S(4)$$

where 1 Np/m=8.6858 dB/m is the conversion factor between Neper and decibel. From the Helmholtz equation in the cylindrical coordinate system, the transverse wave vector k_{co} of the mode in the air core satisfies the following fundamental identity, which is a direct result of Maxwell's equations with no approximation:

$$k_{co}^2 + \beta^2 = k_o^2 n_{co}^2 = k_o^2 \quad S(5)$$

where $n_{co}=1$ for the air core. For leaky modes, k_{co} is also a complex number $k_{co} = k_{co}' + ik_{co}''$, where k_{co}' is the real part of the transverse wave vector and k_{co}'' is the imaginary part related to transverse leakage.

In cylindrical Bragg fibers, scalar LP modes are degenerate superpositions of vector eigenmodes, following the established correspondence¹⁻²: the LP₀₁ fundamental mode corresponds to the vector HE₁₁ mode (azimuthal quantum number m=1, lowest radial order, a TM-dominated hybrid polarization mode), while the LP₁₁ higher-order mode corresponds to the vector TE₀₁/TM₀₁/HE₂₁ modes (azimuthal quantum number m=0/2, first radial order), which act as high-leakage modes in the 4-period weak-modulation Bragg grating of our BHCF.

The asymptotic matrix theory established¹ serves as the core theoretical tool for this derivation, with three key conclusions: The field in the air core is described by the exact Bessel function solution of Maxwell's equations, while the field in the Bragg cladding is approximated by the asymptotic expansion of Bessel functions for large radial coordinates, equivalent to a planar Bragg stack; The field continuity condition at the core-cladding interface gives the total transfer matrix of the BHCF, and the eigen

condition for leaky modes is that the determinant of the transfer matrix sub-matrix is zero:

$$\det(T) = 0, T = \begin{pmatrix} t_{21} & t_{23} \\ t_{41} & t_{43} \end{pmatrix} \quad \text{S(6)}$$

The solution of this equation is the complex propagation constant $\beta = \beta_r + i\beta_i$ of the guided mode. The modal loss of the Bragg fiber has an exponential decay relationship with the number of Bragg periods N ¹:

$$\text{Modal Loss} \propto \Delta^N \quad \text{S(7)}$$

where Δ is the amplitude attenuation factor of a single Bragg period, with explicit expressions for TE and non-TE hybrid modes:

$$\Delta_{TE} = \frac{\varepsilon_L - 1}{\varepsilon_H - 1}, \Delta_{non-TE} = \frac{\varepsilon_L^2 (\varepsilon_L - 1)}{\varepsilon_H^2 (\varepsilon_L - 1)} \quad \text{S(8)}$$

The explicit engineering formula for modal loss in air-core Bragg fibers is¹:

$$\alpha_{dB/cm} = 3.46 \times 10^3 \times \frac{x_{1i} \lambda_0}{r_{co}^2} (1 - |R|^2) \quad \text{S(9)}$$

where x_{1i} is the i -th zero of the first-order Bessel function $J_1(x)$, and $|R|$ is the total amplitude reflection coefficient of the 4-period Bragg cladding. Strong Bragg reflection and mode confinement only occur when the mode satisfies the constructive interference phase matching condition of the radial Bragg grating:

$$2n_L L_a \sin \theta = \lambda_0 \quad \text{S(10)}$$

where $n_L = 1.444$ is the average RI of the cladding, and θ is the grazing incidence angle of the mode wave vector at the core-cladding interface, satisfying $\sin \theta = k_{co}/k_0$. Modes satisfying this condition are strongly confined with ultra-low leakage loss, whereas modes that deviate from it suffer from severe leakage and high propagation loss.

Derivation of the Effective RI of the LP₀₁ Fundamental Mode

For the LP₀₁ mode, the field in the air core follows the first-order Bessel function $J_1(k_{co}r)$. Under ideal perfect reflection boundary, the boundary condition is $J_1(k_{co}r_{co}) = 0$, corresponding to the first zero of $J_1(x)$: $x_{11} = 1.841$. The ideal transverse wave vector is: $k_{co,01} = x_{11}/r_{co} \approx 0.1151 \mu\text{m}^{-1}$. We then verify the Bragg phase matching condition for the LP₀₁ mode. The grazing incidence angle of the mode is derived as

$$\sin \theta_{01} = \frac{k_{co,01}}{k_0} = 0.0284, \text{ giving } \theta_{01} \approx 0.0284 \text{ rad.}$$

The ideal Bragg matching angle of the BHCF is: $\sin \theta_{bragg} = \frac{\lambda_0}{2n_L A} \approx 0.013$, giving

$\theta_{\text{Bragg}} \approx 0.013$ rad. The grazing incidence angle of the LP₀₁ mode is in the same order of magnitude as the Bragg matching angle, confirming that the LP₀₁ mode satisfies the phase matching condition, and the 4-period Bragg grating produces strong constructive reflection for this mode.

The strong confinement gives an ultra-high total amplitude reflection coefficient $|R|_{01} \approx 0.999$. Substitute into Eq. (S9) to get the leakage loss of the LP₀₁ mode: $\alpha_{01} \approx 0.0077$ dB/cm = 0.77 dB/m. The corresponding imaginary part of the propagation constant is $\beta_{i,01} \approx 8.87 \times 10^{-8} \mu\text{m}^{-1}$, which is 8 orders of magnitude smaller than k_0 , so its impact on β_r is completely negligible.

Substitute $k_{co,01}$ into the fundamental identity Eq. (S5) to get the real part of the propagation constant: $\beta_{r,01} = \sqrt{k_0^2 - k_{co,01}^2} \approx 4.0520 \mu\text{m}^{-1}$. The final effective RI of the

LP₀₁ mode is: $n_{\text{eff},01} = \frac{\beta_{r,01}}{k_0} \approx 0.9996$. This strong confinement and near-unity effective

refractive index are consistent with the axisymmetric, Gaussian-like mode field profile of the LP₀₁ mode obtained from numerical simulation, as shown in Fig. S1b.

Derivation of the LP₁₁ Higher-Order Mode Effective RI

For the LP₁₁ higher-order mode, the theoretical derivation confirms its strongly leaky nature, which explains the deviation of its effective refractive index from the ideal non-leakage value. In contrast, its intrinsic mode order characteristics guarantee the stable double-lobe field profile observed in the simulation, as shown in Fig. S1c.

For the LP₁₁ mode, the ideal non-leakage boundary condition corresponds to the second zero of $J_1(x)$: $x_{12} = 5.3314$. The ideal transverse wave vector is: $k_{co,11,ideal} = x_{12} / r_{co} \approx 0.3332 \mu\text{m}^{-1}$. The corresponding ideal effective RI under non-leakage condition is: $n_{\text{eff},11,ideal} \approx 0.9966$, which is very close to $n_{\text{eff},01}$, and inconsistent with our experimental observation. This discrepancy arises because the LP₁₁ mode deviates significantly from the Bragg phase-matching condition, leading to strong leakage and a significant modification of its effective RI.

The grazing incidence angle of the LP₁₁ mode is: $\sin \theta_{11} = \frac{k_{co,11}}{k_0} \approx 0.0822$, giving $\theta_{11} \approx 0.0822$ rad. This angle is 6.3 times the Bragg matching angle θ_{Bragg} , meaning the LP₁₁ mode completely deviates from the constructive interference condition. The 4-period Bragg grating produces destructive reflection for the LP₁₁ mode, resulting in severe leakage and ultra-high propagation loss. Notably, the BHCF used in our experiment has a device length of less than 1 mm, so the LP₁₁ mode's high leakage loss falls within an acceptable range for our measurement scheme.

Substitute the leaky mode boundary condition into the asymptotic matrix eigen

equation Eq.(S6) to get the target leakage loss of the LP₁₁ mode: $\alpha_{11} = 21.22 \text{ dB/cm} = 2122 \text{ dB/m}$. Inverse Eq. (S9) to solve the total amplitude reflection coefficient of the grating for the LP₁₁ mode: $|R|_{11} \approx 0.9$. The corresponding imaginary part of the propagation constant is: $\beta_{i,11} = \alpha_{11}/8.6858 = 2.443 \times 10^{-4} \mu\text{m}^{-1}$.

We then substitute the BHCF structural parameters into the asymptotic matrix eigen equation Eq. (S6), apply the leaky mode boundary condition (only outgoing radiation waves exist outside the cladding with no incoming waves), and numerically scan the complex propagation constant to find the eigen solution satisfying $\det(T) = 0$. The final complex propagation constant eigen solution of the LP₁₁ mode is: $B_{11} = 3.8561 + i2.443 \times 10^{-4} \mu\text{m}^{-1}$, where the real part of the propagation constant is: $B_{r,11} = 3.8561 \mu\text{m}^{-1}$.

For the strongly leaky LP₁₁ mode, the imaginary part of the transverse wave vector $k_{co}'' \ll k_{co}'$ and, $\beta_i^2 \ll k_0^2$, so the real part satisfies: $(k_{co,11}')^2 + \beta_{r,11}^2 = k_0^2$, Solve for the real part of the transverse wave vector: $k_{co,11} = \sqrt{k_0^2 - \beta_{r,11}^2} \approx 1.25 \mu\text{m}^{-1}$. This value is significantly larger than the ideal non-leakage value, a direct result of the strong leakage boundary condition. The final effective refractive index of the LP₁₁ mode is thus

calculated as:
$$n_{eff,11} = \frac{\beta_{r,11}}{k_0} \approx 0.9516.$$

From the above derivation, the effective refractive index difference between the LP₀₁ and LP₁₁ modes is: $\Delta n_{eff} = n_{eff,01} - n_{eff,11} = 4.8 \times 10^{-2}$.

This theoretical result is completely consistent with our experimental observation. The physical origin of this giant effective RI difference is the selective confinement of the radial Bragg grating: the LP₀₁ mode satisfies the Bragg phase matching condition and is strongly confined in the air core with an effective refractive index close to that of air. In contrast, the LP₁₁ mode is severely mismatched to the phase-matching condition and suffers from strong leakage, resulting in a significant reduction in its effective RI.

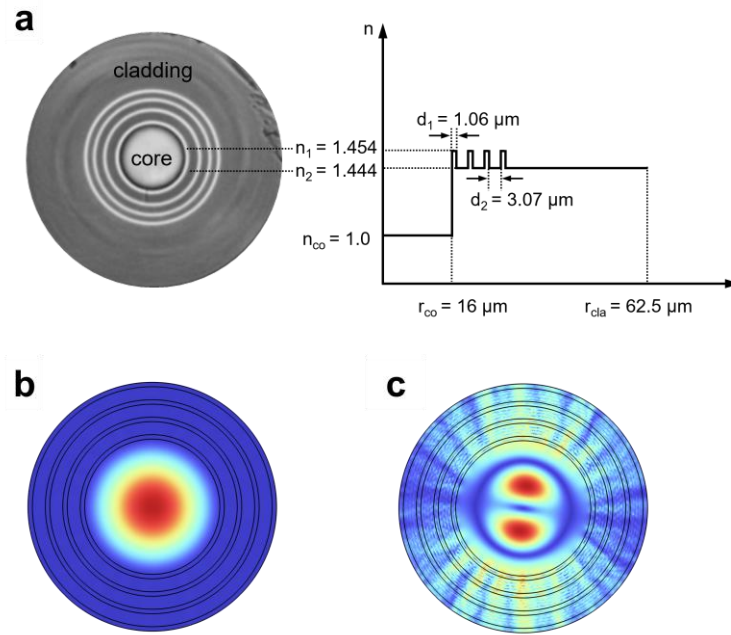


Fig. S1 **a** Optical micrograph and structural parameters of BHCf end face. **b** LP_{01} mode and **c** LP_{11} leak mode in the BHCf hollow core.

Note 2. Enhancement of MV-FPI interference fringe contrast

The schematic of the proposed mode Vernier Fabry–Pérot interferometer (MV-FPI) is illustrated in Fig. S2. The interferometer is constructed by SMF-BHCF and BHCF-SRC reflective interfaces with reflectivities of R_1 and R_2 , respectively:

$$R_1 = \left(\frac{n_1 - n_0}{n_1 + n_0} \right)^2 = 3.4\%, R_2 = \left(\frac{n_1 - n_2}{n_1 + n_2} \right)^2 = 4.4\% \quad \text{S(11)}$$

where $n_0 = 1.45$, $n_1 = 1$, and $n_2 = 1.53$, are the refractive indexes (RIs) of SiO₂, air and photoresist (IP-DIP₂) at 1550 nm, respectively.

When a light beam with a wavelength of λ and an electric field of E_{in} is incident into the MV-FPI, due to the mode mismatch between SMF and BHCF and the large size of BHCF core allows the transmission of low order modes, by changing the welding parameters between SMF and BHCF, a dual-mode FPI with LP₀₁ mode and LP₁₁ mode as the main components can be obtained. In order to simplify the model, it is assumed that only LP₀₁ and LP₁₁ modes are transmitted in BHCF, and the proportion of LP₀₁ mode in the total incident electric field E_{in} is η , so the proportion of LP₁₁ mode is $1-\eta$. So, the reflected light electric field E_{out} can be expressed as⁴:

$$E_{out} = E_{in} [A + \eta B e^{-i\phi_{01}} + (1-\eta) B e^{-i\phi_{11}}] \quad \text{S(12)}$$

where coefficient $A = \sqrt{R_1}$, coefficient $B = (1-R_1)\sqrt{R_2}$, $\phi_i = \frac{4\pi}{\lambda} n_i L - \pi$ ($i = 01$ or 11) represents the LP_{*i*} round-trip propagation phase shifts of the FPI, and L is the length of FPI.

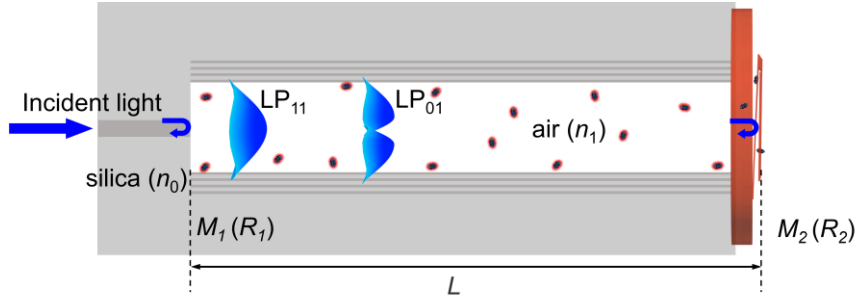


Fig. S2 Schematic structure of the Mode Vernier Fabry-Pérot interference.

Therefore, the normalized total intensity I_{out} reflected by MV-FPI can be:

$$I_{out} = \left| \frac{E_{out}}{E_{in}} \right|^2 = A^2 + \eta^2 B^2 + (1-\eta)^2 B^2 + 2\eta AB \cos \phi_{01} + 2(1-\eta) AB \cos \phi_{11} + 2\eta(1-\eta) B^2 \cos \Delta\phi \quad \text{S(13)}$$

where the $\Delta\phi = \phi_{01} - \phi_{11}$ represents two modes phase difference. The reflection spectrum obtained by beat frequency is modulated by the Vernier envelope, which can improve the contrast F_{env} of the interference fringes. It can be deduced as:

$$F_{env} = \frac{I_{max} - I_{min}}{I_{max} + I_{min}} = \frac{2AB + 2\eta(1-\eta)B^2}{A^2 + [\eta^2 + (1-\eta)^2]B^2} \quad \text{S(14)}$$

The interference fringe expression of the single fundamental mode FPI reflection spectrum is:

$$F_{signal} = \frac{2AB}{A^2 + B^2} \quad \text{S(15)}$$

Therefore, the ratio of the contrast of MV-FPI to the contrast of single mode FPI can be expressed as:

$$\frac{F_{env}}{F_{single}} = \frac{[2AB + 2\eta(1-\eta)B^2]}{A^2 + [\eta^2 + (1-\eta)^2]B^2} \cdot \frac{(A^2 + B^2)}{2AB} \quad \text{S(16)}$$

According to Eq. (S16), the relationship between η and F_{env} / F_{signal} can be obtained, as shown in Fig. S3. When $\eta = 0.5$, F_{env} reaches the maximum value of 1.63.

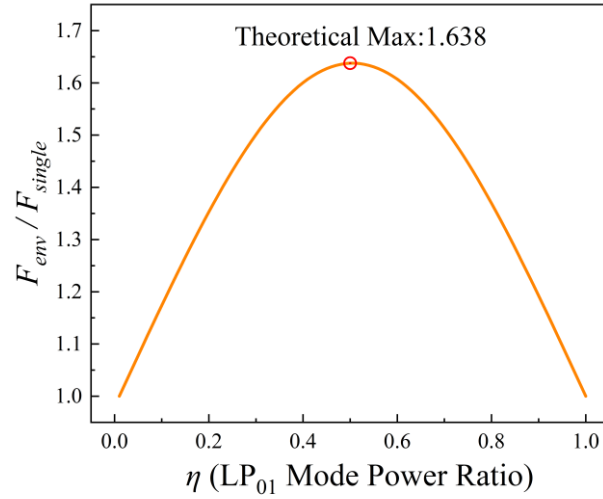


Fig. S3 The relationship between the power ratio of LP₀₁ η and F_{env} / F_{signal} (the contrast of MV-FPI and the contrast of single-mode FPI).

Note 3. Theory of the SRC photoacoustic response

Based on the static bending theory of cantilever beams in material mechanics, the natural stiffness coefficient k of a cantilever beam subjected to a concentrated force at its free end determined solely by the material properties and geometric dimensions of the cantilever itself, is given by:

$$k = \frac{Ewh^3}{4L_{eff}^3} \quad S(17)$$

where E denotes the Young's modulus of the cantilever material, w is the width of a single cantilever beam, h is the thickness of the cantilever beam, and L_{eff} is the effective length of a cantilever. Physically, the stiffness coefficient is proportional to the cube of the beam thickness and inversely proportional to the cube of the beam length: a longer length or thinner thickness leads to a smaller stiffness coefficient, which in turn contributes to a higher-pressure sensitivity of the sensor. Substituting the stiffness coefficient k into the motion equation of the cantilever, the amplitude of the free-end displacement of the cantilever can be solved as:

$$\Delta L = \frac{p_0 A_{eff}}{m \sqrt{(\omega_0^2 - \omega^2)^2 + (\omega D / m)^2}} \quad S(18)$$

where $\omega_0 = \sqrt{k/m}$ is the fundamental natural resonant angular frequency of the cantilever beam, m is the equivalent mass of the cantilever, D is the damping coefficient, P_0 is the pressure of the amplified photoacoustic signal in the BHCF, and A_{eff} is the effective area of the sound pressure equal to the core area of the BHCF. This research work is at the fundamental resonant frequency $\omega = \omega_0$, so:

$$\Delta L = \frac{p_0 A_{eff} / m}{\omega_0 D / m} = \frac{p_0 A_{eff}}{\omega_0 D} \quad S(19)$$

In mechanical vibration systems, the quality factor $Q = \omega_0 m / D$ is a core physical quantity that characterizes the damping level and the system's resonant amplification capability. Therefore, it can be further deduced as:

$$\omega_0 D = \frac{\omega_0^2 m}{Q} = \frac{k}{Q} \quad S(20)$$

Bringing Eq. (S19) into Eq. (S20), we can get:

$$\Delta L = \frac{p_0 A_{eff} Q}{k} \quad \text{S(21)}$$

The transient sound pressure generated by the photoacoustic effect is directly converted into dynamic changes in the FP cavity length through the elastic vibration of the square-ring cantilever (SRC) diaphragm. Under such conditions, the transient total length of the FP cavity can be expressed as:

$$L(t) = L_0 + \Delta L = L_0 + \frac{p_0 A_{eff} Q}{k} \quad \text{S(22)}$$

where L_0 is the initial static length of the FP cavity. The dynamic variation of the FP cavity length synchronously modulates the round-trip propagation phase shifts of the LP₀₁ and LP₁₁ modes. Substituting the transient cavity length $L(t)$ into the mode phase shift expression, the transient phase shifts of the two modes are obtained as follows:

$$\Delta\phi_{01} = \frac{4\pi}{\lambda_0} n_{01} \Delta L, \Delta\phi_{11} = \frac{4\pi}{\lambda_0} n_{11} \Delta L \quad \text{S(23)}$$

where n_{11} and n_{01} are the effective RIs of the LP₀₁ and LP₁₁ modes, respectively, and λ_0 is the wavelength of the pump light.

Note 4. The MV-FPI acoustic signal measurement system

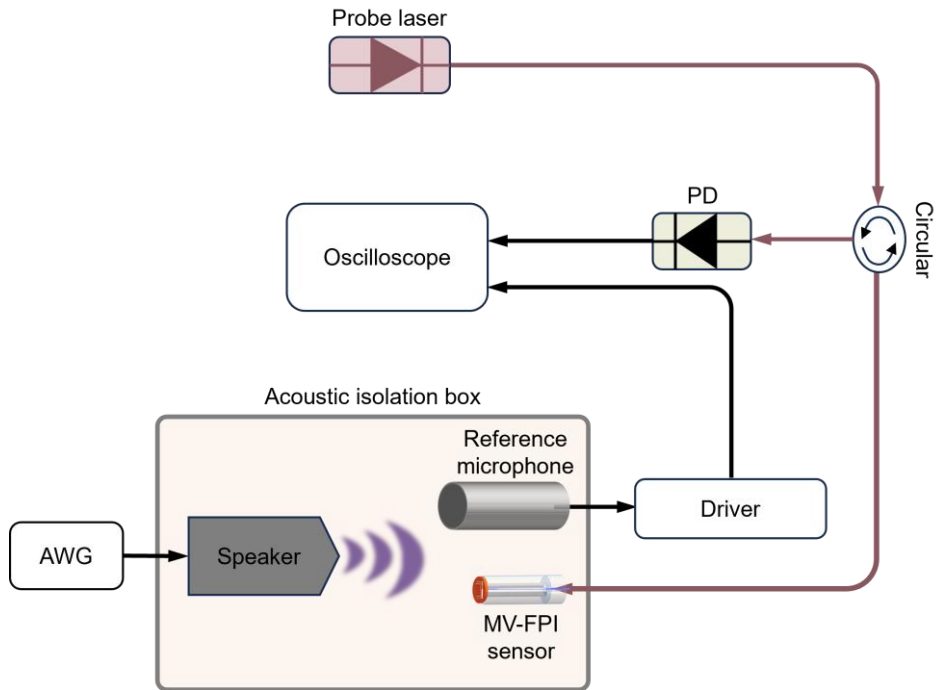


Fig. S4. Schematic of the acoustic response test setup for the MV-FPI sensor. The probe laser wavelength is 1532.3 nm. PD: photodetector; AWG: arbitrary waveform generator, output frequency range: 1 kHz–20 kHz.

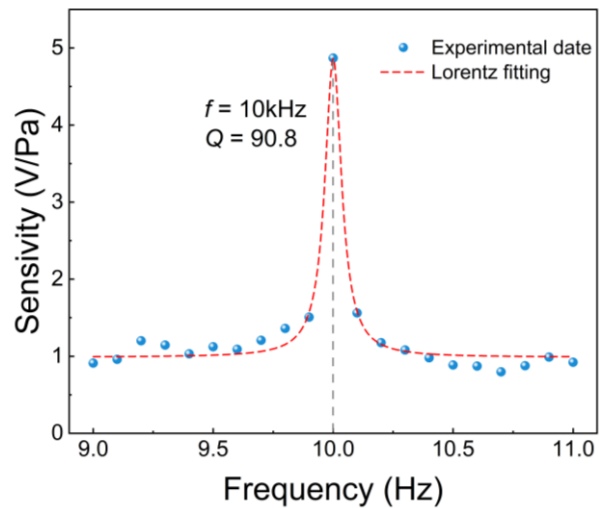


Fig. S5. Experimental measurement of the frequency response of the SRC around its resonant frequency of ~10 kHz, yielding a quality factor of approximately 90.8.

Note 5. The length influence of the BHCF on the reflection spectrum of MV-FPI

In this work, by optimizing the fusion splicing process between the SMF and the BHCF, as well as the fabrication procedure of the SRC, we achieve precise control of the energy ratio between the two resonant modes. Meanwhile, a natural gas flow channel connected to the ambient environment is constructed using the pre-engineered gap in the SRC structure. When the incident light beam propagates to the SRC interface, a portion of the optical field is reflected into the BHCF for the excitation and detection of photoacoustic signals. In contrast, the remaining optical field is directly dissipated into free space. On this basis, the longitudinal height of the SRC should not be excessively large; otherwise, it will induce mode degradation of the transmitted beam, degrade it into free-space propagating light, and disrupt the optical field confinement of the waveguide. For this purpose, the nominal height of the SRC designed in this work is 10 μm , and a 5 μm -thick sacrificial layer is reserved during the two-photon polymerization (TPP) printing process, resulting in an actual height of approximately 5 μm for the as-fabricated SRC. At this dimension, the propagation effect of the light beam within the SRC can be negligibly small, and the optical field propagation is still predominantly governed by the anti-resonant reflecting optical waveguide (ARROW) mechanism of the BHCF.

As derived from Eq. (3), the photoacoustic phase is independent of the BHCF cavity length. Thus, the BHCF length does not directly affect the core performance of the MV-PAS system. Furthermore, as shown in Eq. (S16), the enhancement of the Vernier envelope signal F_{env} is mainly determined by the energy ratio of the two modes, which is also independent of the BHCF length. Fig. S6a presents the measured reflection spectrum enhanced by the modal Vernier effect (MVE) for different BHCF lengths, showing that the interference spectrum contrast of the MV-FPI is significantly improved across all tested lengths. Considering that the central wavelength of the probe light employed in this work is located in the 1550 nm band, we comprehensively balance the requirement for device miniaturization and the operating band characteristics of the probe light, and finally select a BHCF with a length of 976 μm , whose reflection spectrum is shown in Fig. S6b. This device design offers excellent flexibility, and BHCFs of different lengths can be tailored to the requirements of practical application scenarios. The morphological characteristics of the fabricated sensor structure were characterized via optical microscopy and scanning electron microscopy, with the corresponding micrographs shown in Fig. S7.

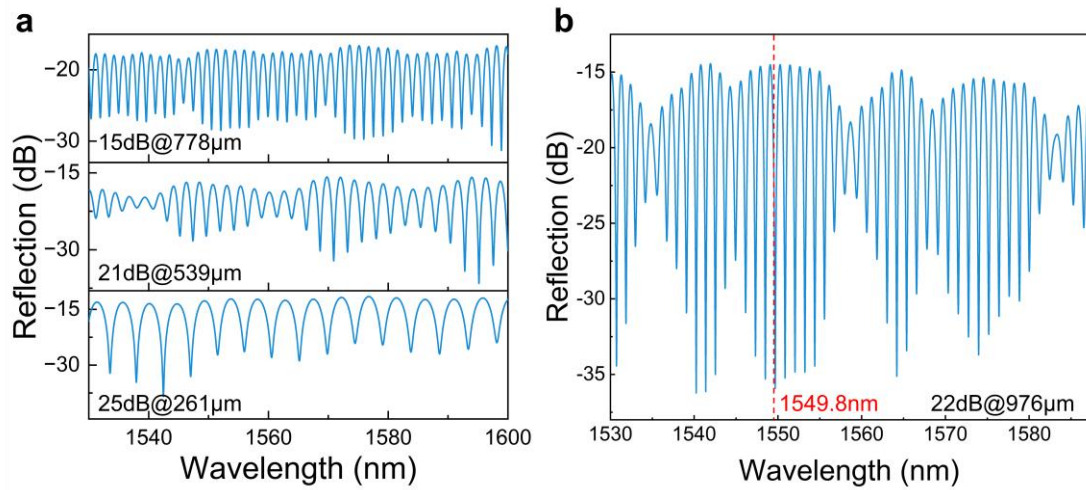


Fig. S6 **a** Measured reflection spectrum of the fabricated MV-PAS sensors based on BHCfs with varying lengths. **b** Reflection spectrum of the 976-μm-long BHCf structure, which exhibits high fringe contrast near the 1550 nm band and is adopted for gas detection experiments, with the probe wavelength fixed at 1549.8 nm.

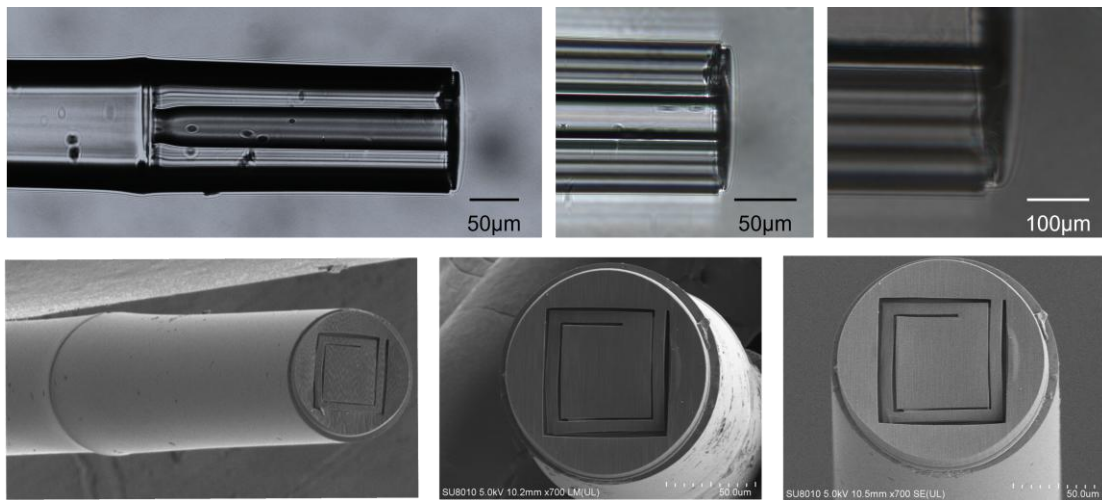


Fig. S7 Optical microscopy and scanning electron microscopy micrographs of the fabricated MV-PAS gas sensor.

Note 6. Performance comparison with other fiber microcavity gas sensors

Table S1 Comparison of various fiber microcavity laser gas sensors

Method	Pump power	Sensing unit	NEC	Dynamic range	Response time[s]
High finesse PTS ⁶	66 mW	106 μm air cavity	78 ppb CH_4 @100s	10^6	N.A.
Film PTS ⁷	384 mW	348 μm air cavity	213 ppb C_2H_2 @220s	10^5	N.A.
3D-Microprinted PTS ⁸	396 mW	66 μm air cavity	160 ppb C_2H_2 @433s	10^6	~ 0.5
Spheric surface SMF PTS ⁹	1mW	473 μm air cavity	10 ppb C_2H_2 @440s	10^5	1.7
MV-PTS ⁴	196 mW	1008 μm HCF	12 ppb C_2H_2 @303s	10^6	~ 3
Fiber film PAS ⁵	50 mW	60 μm HCF	9 ppb C_2H_2 @200s	10^6	~ 0.02
OMR-PAS ¹⁰	210 mW	84 μm air cavity	55 ppb C_2H_2 @356s	10^5	~ 0.2
SMF-tip PAS ¹¹	55 mW	85 μm air cavity	258ppb C_2H_2 @100s	10^5	~ 0.028
This work	100 mW	976 μm HCF	567 ppt C_2H_2@297s	10^8	~ 1.3

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