ARTICLE

Entanglement of propagating optical modes via a mechanical interface

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Many applications of quantum information processing (QIP) require distribution of quantum states in networks, both within and between distant nodes. Optical quantum states are uniquely suited for this purpose, as they propagate with ultralow attenuation and are resilient to ubiquitous thermal noise. Mechanical systems are then envisioned as versatile interfaces between photons and a variety of solid-state QIP platforms. Here, we demonstrate a key step towards this vision, and generate entanglement between two propagating optical modes, by coupling them to the same, cryogenic mechanical system. The entanglement persists at room temperature, where we verify the inseparability of the bipartite state and fully characterize its logarithmic negativity by homodyne tomography. We detect, without any corrections, correlations corresponding to a logarithmic negativity of $E_N = 0.35$. Combined with quantum interfaces between mechanical systems and solid-state qubit processors, this paves the way for mechanical systems enabling long-distance quantum information networking over optical fiber networks.
**Results**

**Theoretical model.** We consider two propagating optical fields (labeled by \( j = A, B \)), from which one can identify a pair of temporal modes with quadratures \( \hat{X}_j, \hat{Y}_j \). We take the variance of these modes to be 1/2 in their ground state. From these modes, one can construct joint Einstein–Podolsky–Rosen (EPR)-type variables \( \hat{X}_\pm = \hat{X}_A \pm \hat{Y}_B \) and \( \hat{Y}_\pm = \hat{Y}_A \pm \hat{X}_B \), which form the basis for various entanglement criteria. We adopt the common Duan–Giedke–Cirac–Zoller (DGCZ) criterion for the inseparability \( \mathcal{I} \), which states that the two modes are inseparable if their variances (\( V \)) satisfy

\[
\mathcal{I} \equiv \frac{V(\hat{X}_\pm) + V(\hat{Y}_\pm)}{2} < 1. \tag{1}
\]

To further quantify this entanglement, one can utilize the system’s covariance matrix, \( \sigma \), which fully characterizes the correlations between various quadratures. From this matrix, it is straightforward to calculate the symplectic eigenvalues of its partial transpose, \( \bar{\sigma} \). These eigenvalues offer a condition for separability (2\( \nu_\perp \geq 1 \)), as well as a tool to calculate a common measure of entanglement, the logarithmic negativity

\[
\mathcal{E} = \log_2 \left( \frac{1}{\mathcal{F}} \right)
\]

which fully characterizes the structure of the soft-clamped mechanical resonator (Si3N4 in white, holes in black). Exiting the cavity, the optical fields possess nonlocal correlations, illustrated by the squeezed phase space ellipses. After the cavity, the two lasers are physically separated and detected simultaneously by balanced homodyne detectors, with local oscillators locked at phases \( \theta_A, \theta_B \). The top of the figure shows a frequency diagram of the relevant optical modes. The two cavity drives are shown in black, with scattered mechanical sidebands of laser A and B shown in blue and red, respectively. The sideband quadrature modes considered in the paper correspond to combinations of both scattered sidebands, as indicated by the blue and red shaded areas.

interaction preserves the Gaussianity of the state. The following equations of motion link the input and output optical fields

\[
\hat{X}_j^{\text{out}}(t) = -\hat{X}_j^{\text{in}}(t), \tag{2a}
\]

\[
\hat{Y}_j^{\text{out}}(t) = -\hat{Y}_j^{\text{in}}(t) - 2\Gamma_{\text{mb}}(t) + 2\Gamma_{\text{mb}}^* P_{\text{in}}(t) + \sum_{i=A,B} 2\sqrt{\Gamma_{\text{mb}}^{\text{in}}(t)} \hat{X}_i^{\text{in}}(t), \tag{2b}
\]

where \( \hat{X}_j^{\text{in}} \) and \( \hat{Y}_j^{\text{in}} \) are input vacuum noise quadratures, \( \chi_m \) is the mechanical susceptibility, \( \Gamma_m \) is the mechanical energy dissipation rate, * indicates convolution, and \( P_{\text{in}} \) is the mechanical thermal noise operator.

In Eq. (2), we see that the quantum amplitude fluctuations of each laser drive the mechanical system, whose motion is then...
imprinted on the optical phase. This is the same mechanism that
drives ponderomotive squeezing of a single laser\textsuperscript{21}, but in this
case there are also cross-correlations between the lasers. More
insight can be had by moving to the joint mode basis (see
Supplementary Note 2), where one finds that the system
decouples into a sum mode undergoing ponderomotive squeez-
ing, and a difference mode which remains dark to all mechanical
dynamics. It is this squeezing of a joint (nonlocal) mode which
results in the “ponderomotive entanglement” we study here. We
note that Eq. (2) also closely mirror those describing four-mode
squeezing based on the Kerr nonlinearity in glass\textsuperscript{22}, in which the
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Furthermore, since the system is stationary, we drop the time
dependence.

In the following model, we consider the limit of long
filter times, addressed by two lasers with wavelength
\( \approx 796 \text{ nm} \). These lasers are orthogonally polarized and populate
the cavity in two different longitudinal modes separated by \( -0.3 \text{ THz} \), with linewidths of \( \kappa_{\text{a}} = 2\pi \times 13.3 \text{ MHz} \) and \( \kappa_{\text{b}} = 2\pi \times 12.6 \text{ MHz} \).

With this setup, we achieve \( \Gamma_{\text{a}} \approx 2\pi \times 1.35 \text{ kHz} \) and
\( \Gamma_{\text{b}} \approx 2\pi \times 0.89 \text{ kHz} \), which easily exceed the thermal deco-
herence rate \( \Gamma_{\text{m}} \approx 2\pi \times 0.20 \text{ kHz} \). We measure the optical

Homodyne detection allows measurement of optical quad-
raturess in a rotated basis defined by the local oscillator phase, \( \theta \).
Filtering the homodyne signal at frequency \( \Omega \) with a mode
function \( h(t) \) yields the sideband quadratures of a particular
temporal mode at time \( t \):

\[
\begin{align*}
\hat{X}_{j}^{\theta}(t) &= \int_{-\infty}^{\infty} ds \cos(\Omega s) h(t-s) \left( \hat{X}^{\text{out}}(s) \cos(\theta) + \hat{Y}^{\text{out}}(s) \sin(\theta) \right), \\
\hat{Y}_{j}^{\theta}(t) &= \hat{X}_{j}^{\theta} + \pi/2(t),
\end{align*}
\]

where \( \theta \) is the homodyne angle. Note that the quadratures \( \hat{X}_{j}^{\theta} \)
available in a homodyne detector contain a pair of sidebands,
symmetric to the carrier, as illustrated at the top of Fig. 1. As
there are in total four optical modes involved, correlations
between such modes are sometimes called four-mode-squeezing
(see Supplementary Note 3), in contrast to the entangled
microwave modes recently analyzed in a heterodyne scheme\textsuperscript{15}.
In the following model, we consider the limit of long filter times,
in which \( h \) effectively selects a single Fourier component\textsuperscript{20}.
Furthermore, since the system is stationary, we drop the time
argument \( t \) and focus on the ensemble statistics of these modes.

Within this model, one can obtain a simple expression for the
DGCZ inseparability criterion (see Supplementary Note 1 for
detail)

\[
\mathcal{I}_{\text{ideal}}^{\text{DGCZ}}(\Theta) = 1 + 8\Gamma_{\text{qua}}|\chi_{\text{m}}(\Omega)|^2 \Gamma_{\text{dec}} (1 + \cos(2\Theta)) - 4\Gamma_{\text{qua}} \Re \left[ \chi_{\text{m}}(\Omega) \right] \sin(2\Theta),
\]

where \( \Theta \equiv (\Theta_{\text{a}} + \Theta_{\text{b}})/2 \). The first term is the contribution from
shot noise at the detectors. The second term is the contribution from mechanical motion, where the total decoherence rate \( \Gamma_{\text{dec}} = 2\Gamma_{\text{qua}} + \Gamma_{\text{m}}(\eta_{\text{b}} + 1/2) \) includes both quantum backaction sources
and thermal motion. The third term corresponds to correlation
between two beams, again in close analogy to ponderomotive
squeezing\textsuperscript{21}. In practice, there is always optical loss, which admits
vacuum noise that degrades the detected correlations. This is
described by a collection efficiency \( \eta_{\text{a}} = \eta_{\text{b}} \equiv \eta < 1 \), with which
the inseparability of the detected optical states becomes

\[
\mathcal{I}_{\Theta} = \eta_{\text{ideal}} - (1 - \eta).
\]

By defining a combined measurement efficiency \( \eta_{\text{meas}} = 2\eta \Gamma_{\text{qua}} \Gamma_{\text{dec}} \) one can show (see Supplementary
Note 1) that the minimum of \( \mathcal{I}_{\Theta} \) is given by \( 1 - \eta_{\text{meas}}/2 \). By
further calculating the full covariance matrix for this toy model
(see Supplementary Note 1), one can show that

\[
\text{min} \left\{ 2\eta / \mathcal{I}_{\Theta} \right\} = \sqrt{1 - \eta_{\text{meas}}},
\]

that is, the system can generate arbitrarily strong entanglement at \( \eta_{\text{meas}} \to 1 \).

\section*{Experimental setup}

In practice, the optical fields become entangled via their shared interaction with a 3.6 mm \( \times \) 3.6 mm \( \times \) 20 nm soft-clamped Si\(_3\)N\(_4\) membrane\textsuperscript{25}. The vibrational mode of
central defect has a frequency of \( \Omega_{\text{m}} = 2\pi \times 1.139 \text{ MHz} \), and a
quality factor \( Q = 1.04 \times 10^9 \) at a temperature of 10 K, which

As illustrated in Fig. 1, the membrane is inserted in the middle of an optical cavity\textsuperscript{26,27},
addressed by two lasers with wavelength
\( \approx 796 \text{ nm} \). These lasers are orthogonally polarized and populate
the cavity in two different longitudinal modes separated by \( -0.3 \text{ THz} \),
with linewidths of \( \kappa_{\text{a}} = 2\pi \times 13.3 \text{ MHz} \) and \( \kappa_{\text{b}} = 2\pi \times 12.6 \text{ MHz} \).

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herence rate \( \Gamma_{\text{m}} \approx 2\pi \times 0.20 \text{ kHz} \). We measure the optical

Inseparability. We now proceed to characterize the variance of
EPR-type variables, as introduced above, to compare with the
DGCZ criterion. We choose a common basis \( \Theta_{\text{a}} = \Theta_{\text{b}} = 0 \) and
measure, in sequence, the combinations \( \{ \hat{X}_{\Theta}^{\text{a}} \hat{Y}_{\Theta}^{\text{b}} \}, \{ \hat{Y}_{\Theta}^{\text{a}} \hat{X}_{\Theta}^{\text{b}} \} \),
and vacuum noise (by blocking the cavity output). Figure 2a, b
shows histograms of the measured quadrature data for the X
and Y quadratures, along with reference lines for vacuum
noise variance in black. Recalling the joint quadrature definitions,
we note that the DGCZ criterion involves the diagonal and anti-
diagonal variances of the X and Y histograms, respectively.

In the figure, we clearly see squeezing in the former, and near-
vacuum variance in the latter—thus already indicating
violation of the DGCZ criterion. Quantitatively, we find
\( \mathcal{I} = 0.83 \pm 0.02 \text{ (stat.)} \pm 0.03 \text{ (syst.)} \). The statistical error comes
from the number of samples used to estimate the EPR variances
and the vacuum noise. The systematic error arises from the
estimations of the vacuum noise variance, due to residual classical
amplitude noise and mismatch in the photodiode responsivities,
at the balanced homodyne detectors (see Supplementary Note 6).

We also notice that the variances in the orthogonal directions are
at the vacuum level. This does not violate the Heisenberg
uncertainty relation, since the pairs of quadratures \( \{ \hat{X}_{\Theta}^{\text{a}}, \hat{Y}_{\Theta}^{\text{b}} \} \)
and vacuum noise (by blocking the cavity output). Figure 2a, b
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diagonal variances of the X and Y histograms, respectively.

Covariance matrix tomography. Having established entangle-
ment, we now quantify it by reconstructing the covariance matrix
Fig. 2 EPR quadrature statistics. a, b 2D histograms of raw X and Y quadrature data, respectively, for Θ ≈ 0. The black dashed circle indicates (2×) the vacuum noise variance, which has a radius of 1/√2 (note the axes’ scale factor of √2). The solid black ellipses indicate (2×) the covariance ellipse of the measured data. The black arrows indicate the diagonal/anti-diagonal variances which are relevant for calculating the DGCZ criterion. c Homodyne angle dependence of joint quadrature variances. The purple (orange) dots are the sum (difference) quadrature \( \hat{X}_A \) (\( \hat{Y}_A \)). The average of these yields the DGCZ inseparability criterion (green points). The measurement ensembles contain \(-10^4\) samples, such that the statistical standard error of the variance estimators is ~2% of the reported values. This contains both the error in the estimation of the EPR variances and the error in the estimations of the vacuum noise variance. The solid line is the theoretical prediction.

Fig. 3 Covariance matrix of the two optical modes. a Measured (black) and predicted (gray) entries of the covariance matrix. The measurement ensembles contain \(-10^4\) samples, such that the standard error of the variance estimators is ~2% of the reported values. b Matrix representation of the measured data, to highlight the location of the significant nonzero entries.

by Gaussian homodyne tomography. By measuring five different pairs of angles \( \{\theta_A, \theta_B\} = \{0, 0\}, \{\pi/2, \pi/2\}, \{0, \pi/2\}, \{\pi/2, 0\}, \{\pi/4, \pi/4\} \), we obtain all necessary intrasystem and intersystem correlations. The reconstructed covariance matrix and theoretical prediction are shown in Fig. 3. From this experimental data, we find a minimum symplectic eigenvalue \( 2\nu_{-} = 0.79 \), corresponding to \( E_N = 0.35 \).

Frequency-dependent entanglement. The previous results refer to a sideband quadrature mode at a particular frequency, \( \Omega \). We now examine how this entanglement varies as we sweep \( \Omega \) near the mechanical resonance, \( \Omega_m \).

(Nota Note 5). Moreover, similar to the measurement in Fig. 3, we can reconstruct the covariance matrix (and corresponding \( \nu_\) for each frequency bin, as shown in Fig. 4b. We see that, as expected, \( 2\nu_{-} \) serves as a lower bound for the inseparability \( I \). Since this bipartite Gaussian state is approximately symmetric, from \( 2\nu_{-} \) we can calculate the entanglement of formation, which is accepted as a proper measure of quantum correlations as a resource\(^1\). Integrating this quantity over a 30 kHz bandwidth yields an entanglement distribution rate of 753 ebits/s.

We emphasize that the optomechanical interaction which generated the entanglement presented above is fundamentally wavelength independent. To illustrate this, we move laser A to \( \sim 912 \) nm, and repeat the measurements of Fig. 4. As shown in Fig. 5, we observe a DGCZ variance below unity and a minimum symplectic eigenvalue \( 2\nu_{-} = 0.92 \) for a mode centered at \( \Omega = 2\pi \times 1.142 \) MHz with bandwidth 915 Hz. The performance is degraded compared to the previous results, due to less efficient light collection at \( \sim 912 \) nm. Nevertheless, these results establish entanglement of two lasers separated by more than 100 nm in wavelength.

Discussion

In conclusion, we have demonstrated quadrature entanglement of two nondegenerate optical beams via their common radiation–pressure interaction with a mechanical resonator.
optical homodyne tomography. The modes of the two laser frequency \( \Omega_b \) are inseparable (green), as a function of center frequency \( a \) and \( c \). The solid lines in a and b are fit to a full model (see Supplementary Note 1). The dark gray line represents the minimum symplectic eigenvalue \( 2\tilde{\nu}_m \) obtained by homodyne tomography. The modes of the two laser fields are entangled whenever \( 2\tilde{\nu}_m < 1 \). c Theory and d measurements of the inseparability \( I(\Theta, \Omega) \). The green dashed line indicates the measurement shown in b. The horizontal axes are referenced to the bare mechanical frequency \( \Omega_m = 2\pi \times 1.139 \text{MHz} \).

While applications in optical or microwave quantum communication are conceivable (as realized with entangled light fields)

presented here, the memory time is ca. 1 ns even for 10 K operation\(^7\), easily exceeding storage time in optical fibers.

In our work, entanglement is preserved from the cryogenic mechanical mediator all the way to the laser beams analyzed in room-temperature homodyne detectors. This enhances the prospects of a general class of hybrid quantum systems\(^{28,29}\) based on mechanical interfaces, which could harness entanglement between solid-state (e.g., spin or charge based) quantum systems, typically operating at low temperatures, and itinerant optical fields.

From a more fundamental perspective, it would be interesting to explore concepts at the interface of quantum measurement and entanglement. For instance, the optomechanical interaction in this work can also be interpreted as a strong quantum measurement of the mechanics. This system should be well-suited for studying the usually hard-to-access system-meter entanglement\(^{26–28}\).

Data availability

Supporting raw data for Figs. 2–5 are available in the UCPH ERDA repository. The remaining data are available from the corresponding author upon request. The repository is https://doi.org/10.17894/ucph.3e60afa5-6377-4488-ab27-850f1e7e6815.

Code availability

The code used in the analysis of the data is available from the corresponding author upon reasonable request.

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Author contributions
J.C., M.R., and D.M. built the set-up and performed the experiments, analyzed the data and, together with A.S., discussed the results and wrote the paper. A.S. supervised the project.

Competing interests
The authors declare no competing interests.

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