We report on an ultrahigh-vacuum compatible optical assembly to enhance the optical power of two separate laser beams with the aim of creating large-scale, two-dimensional optical lattices for use in experiments with ultracold atoms. The assembly consists of an octagon-shaped spacer made from ultra-low-expansion glass, to which we optically contact four fused-silica cavity mirrors, making it highly mechanically and thermally stable. The mirror surfaces are nearly plane-parallel which allows us to create two orthogonal cavity modes with diameters \( \sim 1 \text{ mm} \). Such large mode diameters are desirable to increase the optical lattice homogeneity, but lead to strong angular sensitivities of the overlap between the two cavity modes. We demonstrate a procedure to precisely position each mirror substrate that achieves a mode overlap of \( 99.9 \pm 0.1 \% \). Creating large optical lattices at visible and near infrared wavelengths requires significant power enhancements to overcome limitations in the available laser power. The cavity mirrors have a customized low-loss mirror coating that enhances the power at a set of relevant wavelengths from the visible to the near infrared by up to three orders of magnitude.

### Fabry-Pérot resonators

Fabry-Pérot resonators [1] for visible and near-infrared light are fundamental tools of laser science and atomic, molecular, and optical physics. They are used to provide feedback around a laser gain medium [2], as mechanical length references for optical frequency standards [3], as strongly-coupled interfaces between quantum emitters and light [4], to sense tiny forces such as gravitational waves [5], to enhance the spectroscopic signals from optically thin samples [6], and to create deep optical traps for ultracold atoms [7]. The most important cavity design parameter for each use case is the finesse \( \mathcal{F} \), which determines the cavity’s power enhancement and its frequency selectivity. Decades of technological improvements have enabled \( \mathcal{F} \) up to \( 4 \times 10^5 \) [3, 8]. The second important design parameter is the diameter of the optical mode in the cavity, given by the radius of curvatures \( R_1 \) and \( R_2 \) of the mirrors and the mirror spacing \( L \). Stable Hermite-Gaussian eigenmodes [2] form between the mirrors when the stability parameter product \( 0 \leq g_1 g_2 \leq 1 \), where \( g_i = 1 - \frac{L}{R_i} \). Stable cavities are desirable because they provide clearly distinguishable eigenmodes that allow working with a given laser beam geometry and a well-defined resonance frequency. However, it can be advantageous to work with nearly unstable cavities: even though they are sensitive to mechanical and thermal imperfections, large mode diameters on short cavity lengths can be achieved. Such large modes are essential to create homogeneous optical lattices to trap ultracold atoms for quantum simulators, transportable frequency standards, and atomic sensors.

In this paper, we demonstrate and characterize cavities with...
mode diameters ~1 mm that operate in the near plane-parallel regime with $g_{1,2} \approx 0.995$. For this purpose, we use one flat and one curved mirror with $R = 10.2$ m, spaced by $L = 50$ mm. We present an optical assembly which has no adjustable parts and contains two such cavities whose fundamental modes cross at right angles. The assembly, shown in Fig. 1(a), consists of two pairs of cavity mirrors, optically contacted [9] to an octagon-shaped spacer such that the fundamental modes of both cavities overlap perfectly, are long-term stable, and do not require alignment. When laser light is coupled into each resonator, a two-dimensional optical lattice forms in the region where both modes overlap. The spacer is made from ultra-low-expansion glass and the mirror coatings are ion-beam-sputtered onto fused silica substrates, which ensures vacuum compatibility and a high thermal stability. We choose a moderate finesse $F = 3 \times 10^2 - 5 \times 10^3$ to enable power enhancements $\Lambda = 10^2 - 10^3$ for multiple wavelengths $\lambda$, which keeps the cavity resonances spectrally broad and reduces laser-frequency-noise to amplitude-noise conversion. As shown in Fig. 1(b) our choice of a moderate finesse and near plane-parallel mirror geometry distinguishes us from other cavity applications, such as in cavity quantum electrodynamics [10, 11], reference cavities [3], cavity interferometers [12], or gravitational wave detectors [5]. In this cavity regime, the overlap between the two orthogonal cavity modes becomes sensitive to two main systematic effects.

First, the optical axis of each resonator is shifted when both cavity mirrors are tilted with respect to each other by a relative angle $\beta$, as illustrated in Fig. 1(c). The relative shift of each optical axis dominates over all other systematic effects in the near plane-parallel regime, because each optical axis shifts according to

$$\Delta h = R \sin \beta,$$

which is completely independent of $L$ [13]. Our goal of achieving large $1/e^2$ mode diameters

$$2w = 2 \sqrt{\frac{n^2}{2}} (L/R(1 - L/R))^{1/4} \propto R^{1/4},$$

thus must be balanced against the technical difficulty of polishing glass spacers with extreme requirements on surface parallelism.

Second, the mode separation of higher-order transverse electric-field modes (TEM) becomes comparable to the linewidth of the cavity resonances, such that we approach a regime of “co-existing modes,” as illustrated in Fig. 1(d). Here, the optical axis of each cavity is unaffected, but a misalignment $\Delta \chi$ of the input beam leads to a shift $\Delta \chi \equiv \epsilon \Delta \chi_c$ of the light circulating in the cavity. Even when the input beam is frequency-stabilized to the fundamental TEM$_{00}$ resonance, a fraction of the input power given by the ratio of transverse mode splitting to the cavity linewidth is coupled into the TEM$_{01}$ mode [13]. If we require a maximal misalignment suppression factor $\epsilon$, we find a lower bound on the required cavity finesse (see Appendix A)

$$F \geq \frac{\pi}{2} \arccos \sqrt{\frac{1}{\epsilon^2 - 1}},$$

which we show in Fig. 1(b) for $\epsilon = 5\%$ as a new bound in the stability diagram. For this $\epsilon$ and a mode-matching to the TEM$_{00}$ mode of 99%, our cavity design suppresses $\Delta \chi_c < 1 \mu m$. From the perspective of ray optics we can understand this effect as the minimum number of cavity round trips required to force the average ray offset to lie on the optical axis within a certain margin of error.

Our cavity spacer surfaces have a parallelism $\leq 1$ arcsec to suppress $\Delta h$ to ~ 49 $\mu m$/arcsec. The octagon-shaped spacer is 50 mm wide and 15 mm high, which allows attaching mirror substrates with 12.7 mm diameters to its sides. These sides have a 4 mm wide bore to allow the formation of cavity modes. Further optical access is enabled by 5 mm-wide bores in the remaining sides of the octagon, and a 20 mm-wide bore in its center allows imaging the optical lattices with a high-numerical-aperture microscope objective.

Since the curved cavity mirrors are attached by optical contacting, they require a flat annulus polished onto the curved substrate. For these substrates, the annulus is a 2 mm-wide flat region around a shallow dimple that is less than 1 $\mu m$ deep, as shown in the surface profile in Fig. 2(a). An annulus with an inhomogeneous width would result in an effective tilt $\beta$ of the curved region, which would lead to additional optical axis shifts as discussed above. For our mirror substrates, this tilt (or wedge error) is smaller than 2 arcsec. The combination of shallow mirror surface, small wedge error, and superpolished surface specifications places extreme requirements on the fabrication of the curved mirror substrates. An ion-beam-sputtered high-reflectivity coating was applied to the mirror substrates after polishing. The multilayered coating is 5 $\mu m$ thick and thus protrudes from the dimple in the substrate. The mirror coating was applied only to the central 3.5 mm of the substrate, such that the masked coating fits the spacer’s 4 mm diameter bores and leaves enough room to position the mirror. Furthermore, the uncoated annulus allows a high bonding strength when the mirror is attached to the spacer. The substrates are 6.35 mm thick, which reduces deformation of the substrate due to mechanical stress induced by the coating.

**Fig. 2.** (a) Surface profile of a curved mirror substrate with a polished annulus. The cross denotes the deepest point of the curved area, and the coating is applied within the circled area. (b) Final step of the building process: after all mirrors except one curved mirror substrate are attached, the optical axis of the first cavity mode is measured. The position of the second curved mirror then fixes the optical axis of the second resonator and determines the overlap between both resonators. An interferometer and a transmission measurement performed with a slit are used to find the optimal overlap. Panels (c) and (d) show typical interferograms to determine the alignment of the substrate annulus with respect to the top surface of the spacer.
Since the final assembly has no adjustable parts, we need to ensure that the two orthogonal cavity modes are well overlapped by precisely positioning the cavity mirrors. We optically contact the mirrors to the spacer in a clean-room environment [14] after having cleaned all glass parts using methods developed for semiconductor fabrication [9]. The spacer together with a mirror is mounted on a stage that fixes the spacer while the mirror to be contacted can be tilted and translated with respect to the surface, as illustrated in Fig. 2(b). We use an imaging system and a camera to position the coated region of the mirror with respect to the spacer bore. The mirror is then pressed onto the spacer to initiate the optical bond [9]. With this setup, we measured a lateral positioning uncertainty between initial and contacted mirror position of 30(5) µm. Using these methods, we attach the first mirror pair, which fixes the orientation and position of the first cavity mode. After that, the second flat mirror is attached to the spacer, and consequently the final curved mirror determines the optical axis of the second cavity and thereby the overlap between the two cavity modes.

Before we contact the final mirror, we measure the mode overlap in situ. For this purpose, we recover an interferogram between the lower mirror surface and the spacer top surface, as shown in Fig. 2(b). The interferogram allows us to monitor the relative angle \( \beta \) between the mirror annulus and the spacer surface. A relative tilt leads to straight fringes in the outer regions of the interferogram and a minimal \( \beta \) results in a mostly homogeneous fringe as depicted in Figs. 2(c) and (d). The angular and lateral alignment is limited by optical resolution: a residual fifth of a fringe corresponds to a residual \( \beta = 1 \) arcsec.

After prealigning the final mirror in this way, we measure the mode overlap with a slit that clips parts of both cavity modes. For the overlap measurement, the input beam is frequency-modulated over a free-spectral range, while the slit is positioned using a translation stage. During this process, we measure the transmission of the TEM\(_{00}\) mode for both cavities with separate photodetectors, as shown in Fig. 3(a). To understand the transmission signals, we can think of cutting the cavity mode with one side of the slit as cutting a Gaussian beam with a knife edge which would yield an error function. Because the slit has two edges, the cavity transmission signal varies similar to a sum of two error functions distributed symmetrically around the mode center. In contrast to a traditional knife-edge measurement, however, the cavity mode can diffract and circulate around a partial obstruction such that 50% transmission does not indicate the edge being at the center of the beam profile. For this reason, we choose a wide enough slit width to observe a clear transmission peak (comparable to the \( 1/\epsilon^2 \) mode diameter), while still obstructing the wings of the mode enough to result in a clear maximum in the TEM\(_{00}\) transmission. In Fig. 3(b), we fit a parabola to all points above 50% of the maximum peak transmission to determine the mode center. Fitting both mode centers yields the displacement \( d \) of both modes orthogonal to the plane spanned by the modes. Once we have minimized \( d \), the final cavity mirror is contacted to the spacer. If necessary, we detach the mirror using a razor blade and repeat the contacting procedure. For the final assembly, we characterize and quantify the overlap with eight independent measurements of \( d \), by rotating and flipping the spacer in its mount with respect to the slit. After weighted averaging, we find \( d = 1(5) \) µm. We quantify the overlap by comparing \( d \) to the \( 1/\epsilon^2 \) beam radius \( w \) at the slit position. We determine a final mode overlap \( 1 - d/w = 99.9 \pm 1.3 \% \).

The overlap measurement of the cavity modes is affected by input beam misalignment according to Fig. 1(c). The slit artificially reduces the finesse by obstructing the wings of each mode and thus we can enter the regime of coexisting modes. We measure a reduction in the cavity finesse due to the presence of the slit from 533(19) to 167(38) at a wavelength of 689 nm. To determine the significance of this systematic effect, we vary the input beam position \( \Delta x \), which we set to zero at the position where we have maximum mode matching to the TEM\(_{00}\) mode (>99%), and measure the transmission maximum of the mode \( \Delta x \), which we set to zero at \( \Delta x = 0 \). In Fig. 3(c), we show \( \Delta x \) versus \( \Delta x \), from which we determine the mode displacement suppression factor \( \epsilon \) with a linear fit. For a slit with a width of 800 µm and a mode with \( w = 396 \) µm at 689 nm, we find \( \epsilon = 0.68(1) \) and \( -0.02(1) \), when \( \Delta x \) is altered along and orthogonal to the slit measurement direction, respectively. In Fig. 3(d), we show how \( \epsilon \) decreases with slit size and verify that \( \epsilon \) decreases dramatically for mirrors with smaller radii of curvature with the same mirror coating. To suppress the sensitivity to input beam misalignments and to achieve a valid characterization of the overlap, we maximize the mode matching to the TEM\(_{00}\) mode.
The assembly contains no moving parts and no glue, with the sidebands generated by an electro-optic modulator. We determine visible mode peaks on a photodetector within one free-spectral range, we measure a mode matching of 99% and an input beam misalignment of 14 µm. Besides the overlap, we characterize the power enhancement factor Λ for each wavelength of interest, using a method that can be applied when the finished assembly has been placed inside a vacuum chamber. We determine Λ by estimating the circulating power $P_c$ with the setup shown in Fig. 4(a). We first measure the cavity finesse from a transmission spectrum, which is calibrated with the sidebands generated by an electro-optic modulator. We then determine the amplitude reflection coefficient $r$ of a single mirror via $r = \pi/(2F) - \sqrt{\pi^2/(4F^2)} + 1$ [2]. The enhancement factor is $\Lambda = P_c/P_i$, where $t$ is the amplitude transmission coefficient of the mirror, which is accessible through the on-resonant cavity transmission coefficient $T_e = P_o/P_i = t^2/(1 - r^2)^2$ [2], measured while a laser with wavelength $\lambda$ and power $P_i$ is locked to a cavity arm via a Pound-Drever-Hall lock [15]. We find $\Lambda = \sqrt{T_e}/(1 - r^2)$. As shown in Fig. 4(b), we find no evidence of nonlinear effects up to circulating powers of 80 W, and measure $\Lambda(689 \text{ nm}) = 147(5)$, $\Lambda(813 \text{ nm}) = 161(5)$, and $\Lambda(914 \text{ nm}) = 1132(13)$.

In conclusion, we have demonstrated an optical assembly containing two orthogonal optical resonators whose optical axes cross in the center of the assembly. Each resonator supports a fundamental mode with a $1/\sqrt{2}$ mode diameter approaching 1 mm. The assembly contains no moving parts and no glue, which is essential to work with cavities this far in the near plane-parallel regime due to strong angular sensitivities. This compact optical assembly provides an optimal mode overlap, does not require realignment, is compatible with operation in ultrahigh vacuum, has excellent mechanical and thermal stability, and enhances the optical power for multiple optical wavelengths of interest. By coupling light into the optical resonators, we can create very large and stable optical lattices which benefit quantum simulators, atomic sensors, and optical lattice clocks by providing a large number of homogeneous microtraps for ultracold atoms. This increase in available lattice sites allows using more atoms to directly improve the signal-to-noise ratio in all these applications. The lattice alignment is insensitive even to strong accelerations, which makes our assembly an ideal candidate to integrate into transportable sensors and clocks for geodesy or space missions [16–19].

**Funding.** This work was supported by funding from the European Union H2020 (UQUAM grant no. 319278 and PASQuanS grant no. 817462). A. J. P. was supported by a fellowship from the Natural Sciences and Engineering Research Council of Canada NSERC (funding ref. no. 517029), and N. Š. was supported by a Marie Skłodowska-Curie individual fellowship (grant no. 844161).

**Disclosures.** The authors declare no conflicts of interest. The cavity assembly is subject of an international patent application (PCT/EP2019/066247).

The current address of N. Š is Institute of Physics, Bijenička cesta 46, 10000 Zagreb, Croatia.

**A. DERIVATION OF EQN. (3)**

Assume that an input beam that would be perfectly mode-matched to the TEM$_{00}$ mode of a cavity with mode waist $w$ is instead displaced by $\Delta x_i \ll w$. The electric field $E_i$ of this displaced input beam can be written as a superposition of the cavity’s TEM$_{00}$ and TEM$_{01}$ mode functions $U_0$ and $U_1$ as [13]

$$E_i(x) = A U_0(x - \Delta x_i) \approx A \left[ U_0(x) + \frac{\Delta x_i}{w} U_1(x) \right].$$

Assuming that the input beam is frequency-stabilized to the transmission peak of the TEM$_{00}$ mode, we estimate the electric field $E_c$ circulating in the cavity by taking into account that the TEM$_{00}$ mode is coupled off-resonantly, given by the transverse mode separation frequency [2]

$$v_{sep} = \frac{\Delta \nu_{FSR}}{\pi} \arccos \left(\sqrt{\frac{3g^2}{2}}\right),$$

where $\nu_{FSR} = c/2L$, and $c$ is the speed of light. The contribution of the $U_1$ mode function to $E_c$ compared to its contribution to $E_i$ is thus attenuated by $\varepsilon \equiv [1 + 4(\Delta v_{sep}/\Delta v)^2]^{-1/2}$. For the circulating field, we find

$$E_c(x) \approx A \left[ U_0(x) + \frac{\varepsilon \Delta x_i}{w} U_1(x) \right],$$

from which we conclude that the displacement of the circulating field is suppressed with respect to the displacement of the input field $\Delta E_c \equiv \varepsilon \Delta x_i$. Combining the definition of the finesse $F = \Delta \nu_{FSR}/\Delta v$ with Eqn. 5, we find Eqn. 3.

**REFERENCES**

1. M. Vaughan, *The Fabry-Perot Interferometer: History, Theory, Practice and Applications* (Taylor & Francis, 1989).

2. B. E. A. Saleh and M. C. Teich, *Fundamentals of photonics* (Wiley, 2007), 2nd ed.

3. W. Zhang, J. M. Robinson, L. Sonderhouse, E. Oelker, C. Benko, J. L. Hall, T. Legero, D. G.Matei, F. Riehle, U. Sterr, and J. Ye, Phys. Rev. Lett. 119, 243601 (2017).

4. S. Dutra, *Cavity Quantum Electrodynamics: The Strange Theory of Light in a Box* (Wiley, 2005).

5. J. Aasi, B. Abbott, R. Abbott, M. Abernathy, K. Ackley, C. Adams, T. Adams, P. Addesso, R. Adhikari et al., Cl. Quantum Gravity 32, 074001 (2015).

6. G. Berden and R. Engeln, *Cavity Ring-Down Spectroscopy: Techniques and Applications* (Wiley, 2010).

7. A. Mosk, S. Jochim, H. Moritz, T. Elsaesser, M. Weidemüller, and R. Grimm, Opt. Lett. 26, 1837 (2001).

8. G. Harry, T. Bodiyu, and R. DeSalvo, *Optical Coatings and Thermal Noise in Precision Measurement* (Cambridge University Press, 2012).

9. Q.-Y. Tong and U. Gösele, *Semiconductor Wafer Bonding: Science and Technology* (Wiley, 1998).

10. A. Kawasaki, B. Braverman, E. Pedrozo-Peñafiel, C. Shu, S. Colombo, Z. Li, O. Özeli, W. Chen, L. Salvi, A. Heinz, D. Levonian, D. Akamatsu, Y. Xiao, and V. Vuletić, *Phys. Rev. A* 99, 013437 (2019).
11. F. Brennecke, T. Donner, S. Ritter, T. Bourdel, M. Köhl, and T. Esslinger, Nature 450, 268 (2007).
12. P. Hamilton, M. Jaffe, J. M. Brown, L. Maisenbacher, B. Estey, and H. Müller, Phys. Rev. Lett. 114, 100405 (2015).
13. D. Z. Anderson, Appl. Opt. 23, 2944 (1984).
14. A. Heinz, “Ultracold strontium in state-dependent optical lattices,” Ph.D. thesis, Ludwig-Maximilians-Universität München (2020).
15. E. Black, Am. J. Phys. 69, 79 (2001).
16. T. E. Mehlstäubler, G. Grosche, C. Lisdat, P. O. Schmidt, and H. Denker, Rep. Prog. Phys. 81, 064401 (2018).
17. S. Origlia, M. S. Pramod, S. Schiller, Y. Singh, K. Bongs, R. Schwarz, A. Al-Masoudi, S. Dörscher, S. Herbers, S. Häfner, U. Sterr, and C. Lisdat, Phys. Rev. A 98, 053443 (2018).
18. S. B. Koller, J. Grotti, S. Vogt, A. Al-Masoudi, S. Dörscher, S. Häfner, U. Sterr, and C. Lisdat, Phys. Rev. Lett. 118, 073601 (2017).
19. J. Grotti, S. Koller, S. Vogt, S. Häfner, U. Sterr, C. Lisdat, H. Denker, C. Voigt, L. Timmen, A. Rolland et al., Nat. Phys. 14, 437 (2018).