Electro-optic interface for ultrasensitive intracavity electric field measurements at microwave and terahertz frequencies: supplementary material

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Table S1. The most important symbols used in this paper.

| Symbol | Meaning |
|--------|---------|
| $\omega$ | probe frequency |
| $\kappa_p$ | total probe rate, a sum of the input coupling losses $\kappa_{ex,p}$ and intrinsic losses $\kappa_{0,p}$ |
| $\epsilon_p$ | dielectric permittivity at probe frequency $\epsilon_p = \epsilon_0 n_{mat}^2$ |
| $A_p$ | effective cross-section of the probe mode in the plasmonic waveguide |
| $\Omega$ | THz frequency |
| $\kappa_{THz}$ | total THz rate, a sum of the input coupling losses $\kappa_{ex,THz}$ and intrinsic losses $\kappa_{0,THz}$ |
| $\epsilon_{THz}$ | dielectric permittivity at THz frequency |
| $A_{THz}$ | effective cross-section of the THz mode in the plasmonic waveguide, $V_{THz} \sim A_{THz} l_{gap}$ |
| $g_{eo}$ | single photon electro-optic coupling rate |
| $g$ | light-enhanced electro-optic coupling rate |
| $C_0$ | single photon cooperativity |
| $C$ | multiphoton cooperativity $C = C_0 n_p$ |
| $n_{th}$ | average number of thermal photons stored in the terahertz antenna cavity |
| $n_p$ | total number of probe photons |
| $P_{out}$ | average output laser power |
| $r_{33}$ | electro-optic coefficient of non-linear medium, with $\chi^{(2)}_{33} = -\frac{1}{2} n_{mat}^4 r_{33}$ |
| $\Gamma_c$ | overlap factor between the THz mode and the probing pulse |
| $l_{gap}$ | length of the plasmonic gap, equals interaction length |
| $t_{int}$ | interaction time $t_{int} = \frac{l_{gap} n_g}{v_0}$ |
| $n_g$ | group index of the femtosecond probe in the plasmonic gap |

1. OPTICAL SETUP

The detection setup used in this experiment is shown in Fig. S1a and resembles largely the one used in ref. [1]. In short, we perform the measurements in a dual wavelength THz time-domain spectroscopy setup, that uses probing pulses at a center wavelength of 1556 nm and of duration of 150 fs. Phase-locked THz pulses are generated by pumping a low temperature grown gallium arsenide photoconductive emitter with pulses that are centered around 778 nm. They have a duration of 150 fs and have been produced by frequency doubling the fundamental pulses of the same oscillator. The probing pulses are coupled via a fiber to the chip by grating couplers. The spectrum of the outcoupled pulses after they propagated through the chip was recorded by an optical spectrum analyser and is shown in Fig. S1b. Frequency components contained in the input pulse are efficiently coupled to the chip over a bandwidth >60 nm. The generation of THz pulses is additionally modulated at a frequency of 1.5 kHz by applying a rectangular bias voltage of 12 V to the photoconductive emitter. The temporal evolution of the THz pulses is shown in Fig. S1d, and their spectral amplitudes in Fig. S1e. They travel through an optical path that is ~50 cm long and under ambient atmosphere (not purged) before they reach the antenna detectors. They are collected from the emitter and focussed onto the detector using a pair of parabolic mirrors of 3 inch diameter and focal length of $f = 76.2$ mm. The THz transient is further confined by the THz antenna to the plasmonic gap and modulates here the transmission of the probing pulses through the chip at THz speeds. The THz induced intensity modulation is measured with a photodiode which delivers a voltage proportional to the intensity modulation. A lock-in amplifier demodulates the measured voltage at the frequency with which the emission is modulated. We delay the probe pulses with respect to the THz pulses using a delay stage and reconstruct in this way the entire terahertz waveform that oscillates in the THz cavity.

The power of the pump pulses is 76 mW, when not explicitly stated differently. The power of the probe pulses is indicated for each of the measurements discussed in this article.

In the present case, the multimode probing beam has a temporal extent of ~400 fs at the plasmonic detectors, which is mainly determined by the dispersion introduced by the 1 m long input fiber we use. The silicon waveguide and the plasmonic wavguide have merely no effect on the probe pulse length, as we show in Fig. S1c. The repetition rate of the laser is 90 MHz.

2. ANTENNA CHARACTERISTICS AND SAMPLE DETAILS

In this section, we discuss the performance of the THz antenna structures presented in the main paper. The THz cavity is composed of an antenna and a plasmonic slot waveguide filled with the nonlinear material 3:1 HD-BB-OH/YLD124, as shown in figures Fig. S2a.
Fig. S1. Optical setup characteristics. a, The probing pulses are coupled via fibers to the chip. The generated THz pulses are focused by a parabolic mirror onto the chip after propagating through ambient air after being emitted by a photoconductive antenna. b, Spectrum of probe pulses transmitted through the MZI detector chip. c, Pulse dispersion is mainly introduced by the fiber used to couple into the chip. d, Time-domain evolution of the THz pulse generated by the PCA and measured by a ZnTe crystal of 1 mm length. e, Spectral amplitudes of the emission. PCA: photoconductive antenna, SMF: single-mode fiber, HWP: half wave plate, IR: infrared, fs: femtosecond, ITO: indium tin oxide.

The metallic antenna provides a strong field confinement to the THz vacuum field. To demonstrate this, we perform CST Microwave simulations in the time and frequency domain. In Fig. S2b, we show the simulated energy density for the antenna at 500 GHz. The energy is clearly highly confined in the plasmonic gap. In Fig. S2c we report the cross-section of the electric field distribution for the antenna at 500 GHz, for the THz and near-infrared mode. A good overlap is visible from the simulation, even at such frequency mismatched fields due to the plasmonic confinement.

Details on the individual dimensions of the detectors are given in table S2 using the notation introduced in Fig. S3 together with other characteristic parameters of the antennae such as the average number of thermal photons at the design resonance and the THz loss rate of the fabricated devices. We calculate the THz loss rate $\kappa_{THz}$ from the linewidth of the measured THz field intensity (squared electric field amplitude). To this purpose, we fit Lorenzian lineshapes according to the formula $I(\omega) \sim \frac{1}{(\omega-\Omega)^2 + \left(\frac{\kappa_{THz}/2}{\pi}\right)^2}$ for our two antenna designs at 220 GHz and 500 GHz, for all investigated plasmonic lengths $l_{gap}$. The measured spectra are shown together with the fit in Fig. S3c and d. We note here that the intrinsic losses of the detector due to absorption in the metal and dielectrics constituting our device are negligible compared to the radiative losses of the antenna.

The waveguide structures have been produced by reactive ion etching by starting from a commercial a silicon-on-insulator chip of 220 nm thickness. A 700 nm thick cladding material has been deposited by chemical vapor deposition of SiO$_2$. The metallic antennae have been deposited by electron beam deposition of a Ti/Au (5 nm/150 nm) layer.

data: Table S2. Parameters of the different antenna designs. Geometric dimensions are given together with the average number of thermal photons stored in the THz cavity $n_{th}$, as well as the total measured losses $\kappa_{THz}$ (for the different lengths $l_{gap}$). The 800 GHz antenna does not exhibit a clear resonance, as can be seen from Fig. S3c in the main manuscript. We note that the intrinsic losses $\kappa_{0,THz}$ due to absorption in the dielectric and the gold are negligible compared to the radiative losses of the antenna.

| $\Omega/2\pi$ (GHz) | W (µm) | w (µm) | l (µm) | $w_{gap}$ (nm) | $h_{gap}$ (nm) | $l_{gap}$ (µm) | $n_{th}$ | $\kappa_{THz}$ (THz) |
|---------------------|--------|--------|--------|----------------|----------------|----------------|--------|------------------------|
| 220                 | 30      | 5       | 200    | 150           | 150           | 10, 20, 30 | 7.93   | 2\pi \times 0.18, 0.15, 0.04 |
| 500                 | 30      | 5       | 70     | 150           | 150           | 10, 20, 30 | 12.01  | 2\pi \times 0.27, 0.23, 0.22 |
| 800                 | 30      | 5       | 40     | 150           | 150           | 10, 20, 30 | 7.32   | -                      |
The electric field of a THz wave is measured in this work with sub-cycle temporal resolution via its interaction with a femtosecond probing pulse and the mixing terms effectively introduce a cross-phase modulation of the probe beam. Overall, our experiments are best probing beam in a medium that enables a second-order nonlinear optical mixing. In the following, we derive the most basic equations that describe this process, starting from a general formalism which will be then stepwise applied to our particular case. To ease the readability, we summarise in table S1 the most important symbols used in this paper.

In general, a nonlinear polarisation \( P^{(2)} \) is generated by the mixing of two arbitrary input electric fields which we call \( E_k \). In the following, we derive the Hamiltonian that describes this interaction [2]. We emphasize that the present platform is in certain aspects different from the majority of experimental scenarios discussed in literature. In particular, the near-infrared beam is a pulse which interacts with the THz wave on a sub-cycle temporal and spatial scale. We therefore describe it as a multimode field, as opposed to a single-mode of an optical cavity. The THz wave instead is treated as a cavity mode with a center frequency determined by the THz antenna which has a total loss rate of \( K_{THz} \). The THz wave frequency is in our experiment much smaller than the bandwidth of the probe pulse and the mixing terms effectively introduce a cross-phase modulation of the probe beam. Overall, our experiments are best described by the unresolved sideband regime.

In our design, we exploit the \( \chi^{(2)}_{33} \) component of the nonlinear susceptibility tensor and neglect the rest. Consequently, all fields participating in the interaction are parallel and polarised along the \( z \) axis (defined as the axis across the plasmonic gap) - we omit the subscript \( z \) in the following to ease the notation. The nonlinear interaction energy arises from the nonlinear polarisation

\[
\langle U^{(2)} \rangle = \int dV \int dt \langle \frac{\partial P^{(2)}}{\partial t} \rangle
\]  

(S1)

The interaction Hamiltonian can be derived from the expression of the energy by using the electric field operators in the second quantisation. For the probe pulse we assume a multi-mode field \( \hat{E}_p(t) = i \int d\omega \sqrt{\frac{\hbar \omega}{2\pi V_p}} (u_p(x,y,z)\hat{a}(\omega))e^{-i\omega t} - h.c. \). For the THz field we assume a single-mode \( \hat{E}_{THz}(t) = i \int \frac{\hbar \omega}{2\pi V_{THz}} (u_{THz}(x,y,z)\hat{a}(\Omega))e^{-i\Omega t} - h.c. \). \( u_p(x,y,z) \) and \( u_{THz}(x,y,z) \) are the three-dimensional spatial field distributions that obey the normalisation \( \int dV |u_p(x,y,z)|^2 = V_p, \epsilon_p \) and \( \epsilon_{THz} \) are the permittivities, \( V_p \) and \( V_{THz} \) are the effective mode volumes and \( \hat{a}(\omega) (\hat{a}(\omega)\dagger) \) and \( \hat{a}(\Omega) (\hat{a}(\Omega)\dagger) \) are the annihilation (creation) operators at the two frequencies. We assume without loss of generality that \( u_p(x,y,z) \) and \( \hat{e}_{THz}(x,y,z) \) are real-valued. The total field is therefore \( \hat{E}(t) = \hat{E}_p(t) + \hat{E}_{THz}(t) \), and \( P^{(2)}(t) = \chi^{(2)}_{33} (\hat{E}_p(t) + \hat{E}_{THz}(t))^2 \). We find:

\[
R_I = \int dV \frac{2}{3} \chi^{(2)}_{33} (\hat{E}_p(t) + \hat{E}_{THz}(t))^3.
\]  

(S2)
We define \( g \) with the integral that describes the effective overlap volume \( V \) between the two interacting fields, we find the single photon electro-optic coupling rate between the participating waves as:

\[
\hat{H}_I = 2\epsilon_0 \chi_{33}^{(2)} \hbar \int d\omega \sqrt{\frac{\hbar \omega^2 \Omega}{8 \epsilon_p \epsilon_{THz} V_p V_{THz}}} \int dV |u_p(x,y,z)|^2 u_{THz}(x,y,z) (\hat{a}^\dagger (\omega + \Omega) \hat{a}(\omega) \hat{a}(\Omega) - h.c.)
\]

(S3)

\[
+ \hat{a}^\dagger (\omega) \hat{a}(\omega - \Omega) \hat{a}(\Omega) - h.c.
\]

(S4)

We further consider only the terms which are linear with \( \hat{E}_{THz} \), and neglect all other terms (mixing of the probe only, two photon processes of the THz) and find \( \hat{H}_I = \int dV \chi_{33}^{(2)} E_p(t) E_{THz}(t) \). In addition, the energy conservation rule has to apply. We make use of the fact that all frequency dependent quantities (\( V_p, \epsilon_p, u_p, \omega \)) change slowly as a function of probe frequency within the bandwidth of the pulse. The interaction Hamiltonian can be reduced to:

\[
\hat{H}_I = 2\epsilon_0 \chi_{33}^{(2)} \hbar \int d\omega \sqrt{\frac{\hbar \omega^2 \Omega}{8 \epsilon_p \epsilon_{THz} V_p V_{THz}}} \int dV |u_p(x,y,z)|^2 u_{THz}(x,y,z) (\hat{a}^\dagger (\omega + \Omega) \hat{a}(\omega) \hat{a}(\Omega) - h.c.)
\]

(S3)

\[
+ \hat{a}^\dagger (\omega) \hat{a}(\omega - \Omega) \hat{a}(\Omega) - h.c.
\]

(S4)

We define \( g_{eo}(\omega) \), the electro-optic coupling rate between the participating waves as:

\[
g_{eo}(\omega) = 2\epsilon_0 \chi_{33}^{(2)} \sqrt{\frac{\hbar \omega^2 \Omega}{8 \epsilon_p \epsilon_{THz} V_p V_{THz}}} \int dV |u_p(x,y,z)|^2 u_{THz}(x,y,z)
\]

(S5)

with the integral that describes the effective overlap volume \( V_{overlap} = \int dV |u_p(x,y,z)|^2 u_{THz}(x,y,z) \).

Therefore,

\[
\hat{H}_I = i \int d\omega \hbar g_{eo}(\omega) (\hat{a}^\dagger (\omega + \Omega) \hat{a}(\omega) \hat{a}(\Omega) - h.c.)
\]

(S6)

\[
+ i \int d\omega \hbar g_{eo}(\omega) (\hat{a}^\dagger (\omega) \hat{a}(\omega - \Omega) \hat{a}(\Omega) - h.c.).
\]

(S7)

Using \( \chi_{33}^{(2)} = -\frac{1}{2} n_{mat}^4 r_{33} \) (\( n_{mat} \) is the material index at the probe wavelength, \( r_{33} \) the electro-optic coefficient), and \( \Gamma_c \) the overlap between the two interacting fields, we find the single photon electro-optic coupling rate

\[
g_{eo}(\omega) = \frac{1}{2} n_{mat}^2 r_{33} \sqrt{\frac{\hbar \Omega}{8 \epsilon_p \epsilon_{THz} V_p V_{THz}}} \Gamma_c
\]

(S8)
In the particular case described here, the three-dimensional spatial field profile of the probe beam in the plasmonic can be factorised as follows: \( u_p(x, y, z) = v_p(x, y) w_p(z) \). Similarly, the three-dimensional spatial field profile of the THz beam in the plasmonic can be factorised as follows: \( u_{THz}(x, y, z) = v_{THz}(x, y) w_{THz}(z) \), with \( w_{THz}(z) = 1 \) everywhere in the gap. These approximations are compatible with the field distributions in the gap shown in Fig. S2b and c. Both the THz field as well as the near-infrared probe are well confined to the plasmonic gap. Consequently, the overlap \( \Gamma_c \) can be approximated by the well-defined overlap of the transverse mode profiles:

\[
\Gamma_c = \frac{\int dV |u_p(x, y, z)|^2 u_{THz}(x, y, z)}{\int dV |u_p(x, y, z)|^2 u_{THz}(x, y, z)} = \frac{\int dA |v_p(x, y)|^2 v_{THz}(x, y) \int dz |w_p(z)|^2}{\int dA |v_p(x, y)|^2 \int dz |w_p(z)|^2} = \frac{A_{overlap}}{A_p}
\]

(59)

4. ELECTRO-OPTIC DETECTION

The electro-optic coupling rate depends on the THz mode volume \( V_{THz} \) and on the electro-optic coefficient \( r_{33} \). Simulations of the electric field energy may be used to determine the effective mode volume that influences the vacuum field. In addition, it depends on the overlap of the THz mode with the probe mode.

We can now use the interaction Hamiltonian above to determine how the mode amplitude of the probe pulse evolve due to the non-linear interaction. For this purpose we will compute the quantum field \( \hat{E}_p^{(2)} \) generated by the interaction of the probe pulse with a quantum terahertz field. The probe field is assumed to be classical, hence \( \hat{a}(\omega) = a(\omega) \) and \( E_p(t) = i \int d\omega \sqrt{\frac{\hbar}{2\epsilon_0}} (u_p(x, y, z) a(\omega) e^{-i\omega t} - h.c.) \). Without loss of generality, we assume \( a(\omega) = \sqrt{\frac{\nu}{\pi}} \) to be real-valued. This generated quantum field interferes with the classical probe field, for which we assume no depletion. Consequently, the interaction Hamiltonian can be simplified to:

\[
\hat{H}_I = i \int d\omega a(\omega) \hbar g_{eo}(\omega) (\hat{d}^\dagger (\omega + \Omega) \hat{a}(\Omega) - h.c.)
\]

(10)

\[
+ i \int d\omega a(\omega) \hbar g_{eo}(\omega) (\hat{a}(\omega - \Omega) \hat{a}(\Omega) - h.c.)
\]

(11)

**Fig. S4. Single-mode vs multi-mode mixing process a.** In the case that the probe is in a well-defined single frequency mode, the non-linear interaction leads to a side-band formation that is well resolved. b. In the case that the probe is a broadband femtosecond pulse, the interaction creates unresolved sidebands which effectively introduces a phase delay of the femtosecond probing pulse.

We recognise here a linear combination between the beam-splitter Hamiltonian and a squeezing Hamiltonian, which lead to the formation of side-bands through sum and difference frequency generation (SFG and DFG) as shown in Fig. S4. For sake of simplicity, we derive the evolutions first in the single mode picture and will then add up all generated field to calculate the total probe field after interaction. We start from the equations of motion in the Heisenberg picture:

\[
\frac{d}{dt} \hat{a}(\omega + \Omega) = \frac{i}{\hbar} [\hat{H}_I, \hat{a}(\omega + \Omega)]
\]

(12)

\[
\frac{d}{dt} \hat{a}(\omega - \Omega) = \frac{i}{\hbar} [\hat{H}_I, \hat{a}(\omega - \Omega)]
\]

(13)

and find

\[
\frac{d}{dt} \hat{a}(\omega + \Omega) = a(\omega) g_{eo}(\omega) a(\Omega)
\]

(14)

\[
\frac{d}{dt} \hat{a}(\omega - \Omega) = -a(\omega) g_{eo}(\omega) a^\dagger (\Omega)
\]

(15)

We impose that the interaction time \( t_{int} \) is equal to the propagation time of the probe through the plasmonic gap \( t_{int} = \frac{\lambda_{THz}}{c_0} \). We find that upon interaction, the side-bands are generated with the following amplitudes:

\[
\hat{a}(\omega + \Omega, t_{int}) = a(\omega) g_{eo}(\omega) a(\Omega) t_{int}
\]

(16)

\[
\hat{a}(\omega - \Omega, t_{int}) = -a(\omega) g_{eo}(\omega) a^\dagger (\Omega) t_{int}
\]

(17)
Therefore, after interaction, the resulting electric field of the probe in one arm of the interferometer is:

\[ \hat{E}_{\text{res},1}(t) = E_p(t) + \hat{E}_p^{(2)}(t) \]  

(S18)

Where \( \hat{E}_p^{(2)}(t) = \hat{E}_p^{(2),+}(t) + \hat{E}_p^{(2),-}(t) \) describes the sum of all fields generated in all side-bands by SFG (\( \hat{E}_p^{(2),+}(t) \)) or DFG (\( \hat{E}_p^{(2),-}(t) \)) which obey the following relationships:

\[ \hat{E}_p^{(2),+}(t) = i \int d\omega \sqrt{\frac{\hbar \omega}{2e_p V_p}} u_p(x,y,z) a(\omega + \Omega, t) e^{-i(\omega + \Omega) t} - \text{h.c.} \]  

(S19)

\[ \hat{E}_p^{(2),-}(t) = i \int d\omega \sqrt{\frac{\hbar \omega}{2e_p V_p}} u_p(x,y,z) a(\omega - \Omega, t) e^{-i(\omega - \Omega) t} - \text{h.c.} \]  

(S20)

From here, we combine the equations above and find the total generated field

\[ \hat{E}_p^{(2)}(t) = \int d\omega \sqrt{\frac{\hbar \omega}{2e_p V_p}} u_p(x,y,z) a(\omega) g_{\text{eo}}(\omega) t_{\text{int}}(e^{-i\omega t} a(\Omega) e^{-i\Omega t} - a^\dagger(\Omega) e^{i\Omega t} - \text{h.c.}) \]  

(S21)

We recognise in the formula above the field operator of the THz mode. For simplicity, we define \( \hat{E}_{\text{THz}}^{(1)}(t) = \hat{E}_{\text{THz}}(t) / E_{\text{THz}}^{\text{vac}} = i(a(\Omega) e^{-i\Omega t} - a^\dagger(\Omega) e^{i\Omega t}) \) and define the following phase delay: \( \Delta \phi(t) = \frac{1}{2} n_m^2 q_{33} r_{333} \omega \Gamma_{\text{int}} t_{\text{THz}} \hat{E}_{\text{THz}}(t) \). For the case of \( \langle \Delta \phi(t) \rangle << 1 \), we can approximate \( -i \Delta \phi(t) = e^{-i\Delta \phi(t)} - 1 \) and find that

\[ \hat{E}_p^{(2)}(t) = -E_p(t) + i \int d\omega \sqrt{\frac{\hbar \omega}{2e_p V_p}} u_p(x,y,z) a(\omega) (e^{-i(\omega t + \Delta \phi(t))} - \text{h.c.}) \]  

(S22)

Finally, the resulting field is therefore

\[ \hat{E}_{\text{res},2}(t) = i \int d\omega \sqrt{\frac{\hbar \omega}{2e_p V_p}} u_p(x,y,z) a(\omega) (e^{-i(\omega t + \Delta \phi(t))} - \text{h.c.}) \]  

(S23)

In the second arm, the poling of the organic nonlinear material has been done using the opposite but equal poling voltage, and as such, the total field of the probe after the interaction with the second arm can be retrieved by replacing \( r_{33} \rightarrow -r_{33} \) (and hence \( g_{\text{eo}}(\omega) \rightarrow -g_{\text{eo}}(\omega) \), \( \hat{E}_{p}^{(2)} \rightarrow -\hat{E}_{p}^{(2)} \) or \( \Delta \phi(t) \rightarrow -\Delta \phi(t) \)). Therefore, here the resulting field is \( \hat{E}_{\text{res},2} \) with:

\[ \hat{E}_{\text{res},2}(t) = i \int d\omega \sqrt{\frac{\hbar \omega}{2e_p V_p}} u_p(x,y,z) a(\omega) (e^{-i(\omega t - \Delta \phi(t))} - \text{h.c.}) \]  

(S24)

We recognise from this final formula that the interaction effectively introduces a phase delay of the probe which is linearly dependent on the incident THz quantum field.

**A. Homodyne mixing by the on-chip interferometer**

We compute the intensity modulation introduced by the non-linear mixing in the interferometric scheme realized on-chip. The interferometer is operated at a \( \frac{\pi}{4} \) nominal delay. With the two equations above, we find that the total transmitted intensity is given by

\[ I_{\text{out}}(t) \sim |\hat{E}_{\text{res},1}(t) + \hat{E}_{\text{res},2}(t + \frac{\pi/2}{\omega})|^2 \]  

(S25)

\[ I_{\text{out}}(t) \sim |E_{\text{res},1}|^2 + E_{\text{res},1}(t) E_{\text{res},2}(t + \frac{\pi/2}{\omega}) + E_{\text{res},1}(t)^\dagger E_{\text{res},2}(t + \frac{\pi/2}{\omega}) + |E_{\text{res},2}|^2 \]  

(S26)

\[ I_{\text{out}}(t) \sim 2|E_p|^2 + \Delta I_{\text{out}} \]  

(S27)

We now assume for simplicity that \( a(\omega) = \alpha = \sqrt{\pi} p_s g_{\text{eo}}(\omega) = g_{\text{eo}} \) is constant for all probe frequencies because the relative bandwidth of the probe is limited and we are far away from any resonances. The phase delay now becomes \( \Delta \phi(t) = g_{\text{eo} \text{int}} E_{\text{THz}}^n(t) \). We find that

\[ \Delta I_{\text{out}}(t) \sim 4 \Delta \phi(t) |E_p|^2 \]  

(S28)

In addition, we define \( I_{\text{out}} \sim 2|E_p|^2 \) as the output intensity under no incident terahertz field. We find that:

\[ \frac{\Delta I_{\text{out}}(t)}{I_{\text{out}}} = n_s^2 r_{33} k_0 \Gamma_{\text{gap}} \eta_g i \frac{\hbar \Omega}{2e_p V_{\text{THz}} V_{\text{THz}}} (\beta(\Omega) e^{-i\Omega t} - \alpha(\Omega)^\dagger e^{i\Omega t}) \]  

(S29)

(S30)
The overall intensity modulation is thus linear to the THz vacuum field.

Finally, we define the electro-optic operator $\hat{S}_{eo}(t)$ [3], which, by definition, is:

$$\hat{S}_{eo}(t) = \sqrt{B} i \sum_n \frac{\hbar \Omega}{2 \epsilon_{THz} V_{THz}} ((\hat{a}(\Omega) e^{-i \Omega t} - \hat{a}(\Omega)^\dagger e^{i \Omega t}) (S31)$$

with $\sqrt{B} = n_{med}^2 r_{33} k_0 \Gamma c \tau_{gap} n_g$.

### B. Electro-optic detection of different types of THz radiation

In this section, we discuss the expectation value of an electric field measurement by electro-optic sampling on different types of terahertz radiation. We emphasise that, contrary to continuous measurements, the scheme employed here samples a sub-cycle portion of the terahertz wave with a rate equal to the repetition rate of the probe laser, 90 MHz. The time interval between two consecutive measurements is thus 11 ns, much longer than the measured decoherence time of the THz antennae ($\tau_{THz} = k_{THz}^{-1}$), which are on the order of ∼1-10 ps. Therefore, two subsequent measurements are not correlated and also any back-action induced heating of the THz cavity will have decayed until the next measurement.

First, we will consider the case of a statistical mixture of Fock states $|n\rangle$, such as is the case of thermal radiation. They are described by a density matrix $\hat{\rho} = \sum_n P(n) |n\rangle \langle n|$. The bare vacuum state is a sub-category of such a statistical mixture, where all probabilities are zero besides $P(0)$, which is associated with the vacuum state.

$$\langle \hat{S}_{eo}(t) \rangle_{\hat{\rho}} = \sum_n \langle n| \hat{\rho} \hat{S}_{eo}(t) |n\rangle = 0 \quad (S33)$$

for any type of statistical mixture of Fock states. As such, also $\langle 0| \hat{S}_{eo}(t) |0\rangle = 0$.

Second, we will consider the case of a coherent state of terahertz radiation $|\alpha_{THz}\rangle$. In this case, the THz wave contains $\sqrt{B_{THz}} = \alpha_{THz}$ THz photons.

$$\langle \hat{S}_{eo}(t) |\alpha_{THz}\rangle = \langle \alpha_{THz} \hat{S}_{eo}(t) |\alpha_{THz}\rangle = \sqrt{B} \sum_n \frac{n_{THz} \hbar \Omega}{2 \epsilon_{THz} V_{THz}} \sin \Omega t \quad (S34)$$

The above formula can be interpreted for a classical THz field $E_{THz,q}(t) = \sum_n \frac{n_{THz} \hbar \Omega}{2 \epsilon_{THz} V_{THz}} \sin \Omega t$ in the gap of the plasmonic modulator.

In this condition, the THz wave introduces a time-dependent intensity modulation which is linearly dependent on the local THz electric field in the gap $E_{THz,g}(t)$:

$$\frac{\Delta I_{out}(t)}{I_{out}} = 2 \Delta \phi_{THz}(t) \quad (S35)$$

where $I_{out} = \frac{I_{in}}{e^{-\hat{a}^\dagger \hat{a}}}$ is related to the input probe intensity $I_{in}$ and is considered at $\Delta \phi_{THz}(t) = 0$ and a passive imbalance of the interferometer of $\pi/2$. $\Delta \phi_{THz}(t)$ is the phase delay introduced by the THz wave in one arm via the linear electro-optic effect in the space-time volume in which it overlaps with the probe pulse. In the event of phase-matched detection, the phase delay has been shown above to be equal to:

$$\Delta \phi_{THz}(t) = \frac{1}{2} \sum_n \frac{n_{THz} \hbar \Omega}{2 \epsilon_{THz} V_{THz}} l_{gap} k_0 \Gamma n_g \quad (S36)$$

If we consider phase matching, the phase delay has to be considered for each frequency component of the broadband THz pulse individually as it additionally depends on a frequency-dependent term $\text{sinc}(\frac{\Omega n_{g} l_{gap}}{c_0})$:

$$\Delta \phi_{THz}(\Omega) = \frac{1}{2} \sum_n \frac{n_{THz} \hbar \Omega}{2 \epsilon_{THz} V_{THz}} (\Omega) l_{gap} k_0 \Gamma n_g \text{sinc}(\frac{\Omega n_{g} l_{gap}}{c_0}) \quad (S37)$$

The modulation efficiency is thus $\eta = \frac{\Delta \phi_{pp}}{\Delta \phi_{PP}} = 2 \Delta \phi_{THz}^PP$.

The modulated intensity is detected by a photodiode with a typical bandwidth of 2 kHz, and transformed into a linearly dependent voltage $\partial V = G \Delta P$. $G$ is the gain and $\Delta P$ the modulation in power. Clearly, all equations above have been reducted in units of intensity, however, they take an equivalent form for power, or output voltage.

### C. Finding the experimental $\hat{S}_{eo}$ by probing the interface with a weak THz coherent state

We calculate the experimental single-photon electro-optic coupling rate by probing the system with a weak coherent state of THz radiation $|\alpha_{THz}\rangle$ that contains an average number of photons $n_{THz} = |\alpha_{THz}|^2$. Under this assumption, the electric field amplitude of the coherent state is $\sqrt{n_{THz}} E_{THz,\text{vac}}$. Therefore, the peak-peak phase modulation $\Delta \phi_{THz}^PP$ and the peak-peak intensity modulation $\Delta I_{pp}$ depend on $\hat{S}_{eo}$ as follows:

$$\frac{\Delta I_{pp}}{I_{out}} = 2 \sqrt{B} \frac{n_{THz} \hbar \Omega}{2 \epsilon_{THz} V_{THz}} = 4 \hat{S}_{eo} I_{int} \sqrt{n_{THz}} \quad (S38)$$

$$\hat{S}_{eo} = \frac{\Delta I_{pp}}{4 I_{out} I_{int} \sqrt{n_{THz}}} = \frac{\eta c_0}{4 n_g l_{gap} \sqrt{n_{THz}}} \quad (S39)$$
Finally,

\[ g_{cc} = \frac{\Delta \phi_{THz}^p c_0}{2\sqrt{\pi} \theta_{THz} \text{gap} \eta_g} \]  

(S40)

As can be seen from the formula above, the number of input THz photons is crucial to determine the experimental electro-optic coupling rate. This has been now computed from the spectrum of the input terahertz pulse, assuming a cross-section at the focus of 1 mm diameter and a pulse length of 1.5 ps. The experimentally determined photon number contained in the THz pulse is shown in Fig. S5.

Fig. S5. Number of photons contained in the incident THz pulse. The number of photons contained in the incident THz pulse is computed from the emission spectrum shown in Fig. S1e.

D. Losses and cooperativity

The single photon cooperativity \( C_0 \) and the cooperativity related to the probe photons \( C = \langle n_p \rangle C_0 \) depend on the loss rates of the system. They in essence relate the rate at which the nonlinear mixing leads to the generation of a photon, to the rate at which the participating photons are lost by the structure.

The single photon cooperativity \( C_0 \) is:

\[ C_0 = \frac{4 g_{cc}^2}{\kappa_p \kappa_{THz}} \]  

(S41)

with \( \kappa_p, \kappa_{THz} \) the loss rates of the probe and THz photons which we discuss below.

We can define a total cooperativity of the coherent probe with a single THz photon as

\[ C = \langle n_p \rangle C_0 \]  

(S42)

In our case, the total number of photons in the probe is equal to \( \langle n_p \rangle = \frac{P_{out}}{\hbar \omega f_{rep}} \), with \( P_{out} \) the probe power at the output of the plasmonic detector.

The THz loss rate \( \kappa_{THz} \) are reported for our two antenna designs at 220 GHz and 500 GHz as shown in Fig. S3c and d and in table S2.

Fig. S6. Losses. Propagation loss \( a \) (in dB) in the plasmonic gap is given for two different batches of chips.
The total probe loss rate $\kappa_p = \kappa_{0,p} + \kappa_{ex,p}$ is constituted of the sum between intrinsic and extrinsic probe losses, see Fig. S6a. While there are different approaches in literature on how to define the different contributions to the losses, we choose to define the intrinsic probe losses as all lost photons that are not detected in the output photodiode, as for example the photons absorbed in the plasmonic gap. This rate can be computed from the propagation loss shown in Fig. S6b. For a plasmonic length of 10 $\mu$m and a plasmonic width of 150 nm, the absorption is 3 dB, hence $\kappa_{0,p} = 0.5$. Therefore, $\kappa_{0,p} = -\frac{1}{l_{\text{out}}} \ln(\frac{l_{\text{out}}}{l_{\text{in}}}) = 2\pi \times 0.95$ THz.

The extrinsic probe loss rate is defined as the rate at which the probe photons escape from the interaction region. In the present case, this is inversely proportional to the propagation time through plasmonic gap $\kappa_{ex,p} = \frac{1}{l_{\text{out}}}$. For a 10 $\mu$m long plasmonic gap, the loss rate is $\kappa_{ex,p} = 2\pi \times 1.59$ THz.

With this value, we have that the measured single photon cooperativity for the 220 GHz antenna design is $C_0 = 1.6 \times 10^{-8}$ for 220 GHz, $l_{\text{gap}} = 10 \mu$m. To determine the cooperativity related to the probe, we use the input probe power, which is equivalent to doing a post-selection on the measurement. At $P_{\text{out}} = 10^{-3}$ mW, the total number of probe photons is $\langle n_p \rangle = 122873$ we find the cooperativity $C = 0.002$.

E. Analysis of electro-optic coupling rate and single photon cooperativity

We are providing in the following a numerical analysis of the device parameters’ influence on the electro-optic coupling rate $g_{eo}$ and single photon cooperativity $C_0$. This should serve as a guideline for new designs.

Of particular interest are the plasmonic gap width and length, as these two parameters have the strongest influence on the performance of the detector. In particular, they both have a direct impact on the terahertz mode volume and the losses. For simplicity, we approximate the terahertz mode volume by $V_{THz} = V_{\text{gap}}l_{\text{gap}}h_{\text{gap}}F$, where $F$ is a correction factor determined by numerical simulations. The overlap factor which is a function of the gap width is approximated by $\Gamma_c = 15w_{\text{gap}}^{-1} + 0.25$ as found from numerical simulations. Using eq.(8) we can plot the electro-optic coupling rate for a terahertz frequency $\Omega = 2\pi \times 220$ GHz.

![Fig. S7. Electro-optic coupling rate as a function of the gap width and gap length.](image)

It can be seen that the best coupling rates are achieved for narrow gaps, i.e. providing strongest field enhancement [4] and best overlap [5], in combination with shorter gap lengths leading to smallest terahertz mode volumes. Next, Fig. S8 plots the single photon cooperativity, which includes the geometry dependent losses, as a function of the gap width and length. The losses are computed from the measured loss rate (dB/µm) shown in Fig. S6b. Again, the highest values are found for narrow plasmonic gap and small length, i.e. combination of small mode volumes and lower losses.

F. Equivalence with position measurements of mechanical modes in optomechanics

In quantum mechanics, there is an analogy between measurements of quantum vacuum electric fields and measurements of the zero-point motion of mechanical oscillators. Formally, the analogy stems from the description of both mechanical and electromagnetic fields by a quantum harmonic oscillator. As a consequence, the physical description of electro-optical coupling underlying the measurement of electric fields as we discuss here, can be seen in analogy to the description of an optomechanical coupling that enables position measurements performed in the framework of optomechanics. A comprehensive review can be found for example here [6].

The equivalence can be noticed already in the analogy of the interaction Hamiltonian in the two cases. We have shown above that the interaction Hamiltonian that describes the nonlinear electro-optic interaction in the single mode regime is given by a sum of a beam splitter Hamiltonian and a squeezing Hamiltonian

$$\hat{H}_I = i\sqrt{\frac{\hbar}{m\omega}}g_{eo}(\omega)(\hat{a}^\dagger(\omega + \Omega)\hat{a}(\Omega) - \text{h.c.}) + i\sqrt{\frac{\hbar}{m\omega}}g_{eo}(\omega)(\hat{a}(\omega - \Omega)\hat{a}(\Omega) - \text{h.c.})$$

with

$$g_{eo}(\omega) = \frac{1}{2}h\omega^2\hat{S}\hat{S}_3 (\Gamma_c \frac{2\epsilon_{THz}V_{THz}}{\hbar\Omega})$$

This equivalence can be used to provide a basis for design and optimization of terahertz detectors.
We now proceed to derive the noise due to back-action in first approximation. We impose that the first measurement, which occurs without loss of generality at time $t = 0$, measures only quadrature $\hat{X}^{\text{THz}}$ from the cavity resonance, the interaction Hamiltonian that describes the system can be found e.g. in equation (33) of reference [6] (corresponding to $\hat{\text{potential}}$ to reach the highest optomechanical coupling rate. $\delta \hat{X}$ and $\delta \hat{Y}$ are driven cavity modes that can interexchange quanta, here shown in the rotating frame of the laser frequency of the laser $\omega$ (following ref. [6]).

Typically, the mechanical displacement of a microresonator introduces a phase delay in an evanescently coupled probe. The inherently low resonant frequency of the mechanical mode allows a direct measurement with a continuous-wave probe.

### 5. MEASUREMENT BACK-ACTION AND STANDARD QUANTUM LIMIT

One direct implication of the measurement of quantum electric fields via electro-optic interfaces is the concomitant measurement of the two quadratures of an electromagnetic wave. This can be understood from the temporal evolution of the electric field $\hat{E}^{\text{THz}}(t) = \hat{E}^{\text{vac}}(t) + \hat{E}^{\text{eff}}(t)$, which we consider to be a single mode without any loss of generality. As $[\hat{E}^{\text{THz}}, \hat{E}^{\text{THz}}] = \frac{i}{\hbar}$, a measurement of the electric field will inherently introduce imprecision onto a second measurement of the same quantum state at a later time-point. This effect is particularly important in the specific case where instead of one single measurement, two-pulse correlation measurements are performed with the goal e.g. to retrieve the power spectrum of thermal states, including the vacuum state (as done in ref. [3]). In this scenario, two sub-cycle measurements have to be performed within the decoherence time of the probed radiation (corresponding to $\hat{E}^{\text{THz}}(0)$ and $\hat{E}^{\text{THz}}(\tau)$), and, as such, back-action of the first measurement onto the second one at time $\tau$ later is essential. At the standard quantum limit, the total uncertainty of the measurement is minimized through a fine balance between measurement-based imprecision and imprecision noise due to back-action.

The noise in the first measurement will be determined by the imprecision noise $\hat{s}_{\text{EE}}^{\text{imp}}$, which stems from the shot noise in the probe beam. From formula $\Delta_{\text{EE}}^{\text{imp}} = \sqrt{B_{\text{EE}}^{\text{THz}}} \tau$ derived earlier, we find that the shot noise introduces an apparent noise field (noise equivalent field) $\delta E_{\text{EE}}^{\text{imp}} = \frac{1}{\sqrt{B_{\text{EE}}^{\text{THz}}}} \delta (\Delta_{\text{EE}}^{\text{imp}}) = \frac{1}{\sqrt{B_{\text{EE}}^{\text{THz}}}}$. As such, the power spectrum of the noise field due to imprecision is

$$
\hat{s}_{\text{EE}}^{\text{imp}} = \frac{1}{B_{\text{EE}}^{\text{THz}}} = \frac{|\hat{E}_{\text{EE}}^{\text{THz}}|^2}{4 \omega_{\text{EE}}^{\text{THz}} \Omega^{\text{THz}}}.
$$

We now proceed to derive the noise due to back-action in first approximation. We impose that the first measurement, which occurs without loss of generality at time $t = 0$, measures only quadrature $\hat{X}^{\text{THz}}$. This can be insured by a proper definition of $\hat{X}^{\text{THz}}$ and $\hat{Y}^{\text{THz}}$ to account for the physically irrelevant phase. As an outcome of this first measurement, we have now information about one quadrature, at the expense of noise introduced into the second quadrature, owing to their uncertainty relation $\Delta X^{\text{THz}} \Delta Y^{\text{THz}} \geq \frac{1}{2}$. The back-action noise field can be calculated for small $\Omega \tau$ from:

$$
E_{\text{THz}}^{\text{imp}}(\tau) = E_{\text{THz}}^{\text{vac}} \Delta X^{\text{THz}} \sin(\Omega \tau) \geq \frac{E_{\text{THz}}^{\text{vac}} \Omega \tau}{4 \Delta X_{\text{THz}}} = \frac{|E_{\text{THz}}^{\text{vac}}|^2 \Omega \tau}{4 \delta E_{\text{THz}}^{\text{imp}}} = \frac{|E_{\text{THz}}^{\text{vac}}|^2 \Omega \tau}{4 \sqrt{B_{\text{EE}}^{\text{THz}}}}.
$$

Fig. S8. Single photon cooperativity as a function of the plasmonic gap geometry.
As a result, the back-action noise power is:

\[
S_{EE}^{\text{ba}} = \frac{\left| E_{THz}^{\text{vac}} \right|^2 \Omega^2 t_e^2 a^2 e^2 n_p^2}{4}
\]  

(S49)

Since the total noise is given by the sum of the back-action noise and the imprecision noise \(S_{EE}^{\text{tot}} = S_{EE}^{\text{ba}} + S_{EE}^{\text{imp}}\), there is an optimal value of \(n_p\) at which the total noise is optimised. The value of the minimal total noise marks the standard quantum limit and can be found when \(S_{EE}^{\text{ba}} = S_{EE}^{\text{imp}}\), hence

\[
n_p = \frac{1}{g_e^2 a^2 \Omega \tau}
\]  

(S50)

In the case of the 220 GHz antenna with \(l_{\text{gap}} = 30 \mu m\), we find that we reach the standard quantum limit over a measurement time \(\tau = 0.2\pi / \Omega\) if the probe contains \(\sim 10^6\) photons, which corresponds to an average power of 1 mW. For comparison, in the vacuum field measurements of ref. [3], with the given experimental values of the \(g_e = 2\pi \times 1.6 \text{kHz}\), \(t_{\text{int}} = 32\) ps we find that the standard quantum limit is reached at \(n_p = 10^{13}\), which is experimentally unfeasible. At the standard quantum limit, the total noise is two times higher than the imprecision noise. Much below the standard quantum limit, as the case here, the noise is equal to the imprecision (shot) noise.

6. SHOT-NOISE LIMITED DETECTION

The electro-optic signal depends linearly on the intensity (and therefore total power) of the probe pulse that is detected at the output of the interferometric detector. This signature originates from the detection mechanism that utilises the mixing of the THz wave with the probe pulse. The noise of the electro-optic signal is thus determined by the shot-noise of the probe pulses at the photodiode. The amplitude of the shot noise for a measurement with a single probing pulse is given in units of output voltage \(V\) by:

\[
\sigma V_{sn} = \sqrt{2 rf_{exp} GV / r}
\]  

(S51)

where \(e\) is the elementary charge, \(f_{exp}\) the repetition rate of the laser and \(r = 1\) A/W the responsivity of the InGaAs photodiode and \(G\) the gain of the photodiode in (V/W).

Using this equation, we can now determine the signal-to-noise ratio of electro-optic detection for a given \(E_{THz}\). For simplicity, we assume phase matching along the entire length of the plasmonic gap. This assumption applies for the case discussed in Fig 4 of the main text:

\[
\text{SNR}_{\text{sn single pulse}} = \frac{\Delta V}{\sigma V_{sn}} = \frac{V^n_{\text{mat}} f_{\text{THz}}^2 l_{\text{gap}} k^2 q^2 G V / r}{\sqrt{2 rf_{exp} GV / r}}
\]  

(S52)

Clearly, a high probe power yields a higher signal-to-noise ratio at a given detected THz amplitude. In addition, for high probe powers, the noise power spectral density scales linearly with the probe intensity, as expected for shot noise limited detection, shown in Fig. S9a-b. To maximise thus the probe power from 1 \(\mu\)W to potentially 1 mW is a clear goal to be pursued in the future.

![Fig. S9. Noise behaviour of the detector. a, Dependency of the electro-optic signal as a function of probe power. b, Noise power spectral density as a function of probe power scales linearly. The noise scales therefore like the square root of the number of probe photons at high probe powers. At low probe powers, thermal noise is dominant.](image)

7. BROADBAND DETECTION ARRAY

As a final note, we wish to underline on the incredible richness that such a device architecture opens up by discussing a multi-frequency monolithic detector that contains instead of one single interferometer, four different MZI’s in parallel, each designed to operate on different frequency lines: 220 GHz, 500 GHz, 800 GHz and 1 THz, as shown in Fig. S10a. The frequency response exhibited by this detector is shown in Fig. S10b, together with the three measurements of the individual constituent detectors. The broadband array of MZI clearly shows distinct signatures at the different target frequencies.
**Fig. S10.** Response of an array of four co-integrated MZI modulators in one single detector. a, Microscope picture of the broadband array showing 4 parallel MZIs with different bowtie antennas which are designed with a center frequency of 220 GHz, 500 GHz, 800 GHz and 1 THz. b, Spectral response of the broadband array is compared to the constituent individual antenna responses.

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