Geometric filterless photodetectors for mid-infrared spin light
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Supplementary note S1

In this section, we derive the dependence of photoresponse on Stokes parameters, $S_{0,1,2,3}$: $V_{ph} = R_0 \cdot S_0 + R_1 \cdot S_1 + R_2 \cdot S_2 + R_3 \cdot S_3$, with the $R_{1,2,3,4}$ as the respective responsivities to the four Stoke parameters.

We start from the general expression of the spontaneous photoresponse in non-centrosymmetric materials:

$$ J_i = \beta_{ijk} E_j E_k^* $$

(S1)

where the indices $i, j, k$ represent $x$ or $y$, since our device is planar and the incidence of light is normal to our device. $J_i$ is the current density along the $i$ direction. $\beta_{ijk}$ denotes the photoresponse tensor. $E$ is the complex amplitude of the electric field of the incident light. Because the d.c. photocurrent $J_i$ is real, the right-hand side must also be real. Then, the right-hand side should not change when one takes the complex conjugate, hence $\beta_{ijk} = \beta_{ijk}^*$. The complete expression of the overall photoresponse can be written in the form of a matrix:

$$
\begin{pmatrix}
J_x \\
J_y
\end{pmatrix} =
\begin{pmatrix}
\beta_{xxx} & \beta_{xxy} & \beta_{xyx} & \beta_{xys} \\
\beta_{yxx} & \beta_{yxy} & \beta_{yx} & \beta_{yyx}
\end{pmatrix}
\begin{pmatrix}
E_x E_x^* \\
E_y E_y^* \\
E_x E_y^* \\
E_y E_x^*
\end{pmatrix}
$$

(S2)

Note that the Stokes parameters, $S_{0,1,2,3}$, are related to the amplitude of the electromagnetic wave as:

$$
S_0 = E_x E_x^* + E_y E_y^* \\
S_1 = E_x E_y^* - E_y E_x^* \\
S_2 = E_x E_x^* + E_y E_y^* \\
S_3 = iE_x E_y^* - iE_y E_x^*
$$

(S3)

That is:
Therefore, the photoresponse can also be expressed with the bases of $S_{0,1,2,3}$:

$\begin{pmatrix} S_0 \\ S_1 \\ S_2 \\ S_3 \end{pmatrix} = \begin{pmatrix} 1 & 1 \\ 1 & -1 \\ 1 & 1 \\ i & -i \end{pmatrix} \begin{pmatrix} E_x E_x^* \\ E_y E_y^* \end{pmatrix}$ \hspace{1cm} (S4)

That is:

$\begin{pmatrix} J_x \\ J_y \end{pmatrix} = \begin{pmatrix} \beta_{xx} & \beta_{xy} & \beta_{yx} & \beta_{yy} \\ \beta_{yx} & \beta_{yy} & \beta_{yx} & \beta_{yy} \end{pmatrix} \begin{pmatrix} 1 & 1 \\ 1 & -1 \\ 1 & 1 \\ i & -i \end{pmatrix}^{-1} \begin{pmatrix} S_0 \\ S_1 \\ S_2 \\ S_3 \end{pmatrix}$ \hspace{1cm} (S5)

It is straightforward to find the elements of the photoresponse matrix to be real using the relation $\beta_{ijk} = \beta_{kij}^*$, which makes sense in physics because $J_{x,y}$ and $S_{0,1,2,3}$ are real. Therefore, the measured photocurrents are linearly related to the Stokes parameters. Note that the above derivation will be invalid when the photoresponse is strongly dependent on the incident power, where Eq. S1 does not hold anymore.
Supplementary note S2

The dependence of Stokes parameters on the quarter wave plate (QWP) angle is derived below. Given that the incident light is purely coherent and linearly polarized along the x-axis, the incident polarization state can be described with a normalized Stokes vector:

\[ S_{in} = \begin{pmatrix} S_0 \\ S_1 \\ S_2 \\ S_3 \end{pmatrix} = \begin{pmatrix} 1 \\ 1 \\ 0 \\ 0 \end{pmatrix} \]  \hspace{1cm} (S7)

After passing through the QWP, the new polarization state can be conveniently obtained by Mueller calculus. The Mueller matrix of a QWP is:

\[
M_{QWP} = \begin{pmatrix}
1 & 0 & 0 & 0 \\
0 & \cos^2(\varphi) & \sin(2\varphi)\cos(2\varphi) & \sin(2\varphi) \\
0 & \sin(2\varphi)\cos(2\varphi) & \sin^2(2\varphi) & -\cos(2\varphi) \\
0 & \sin(2\varphi) & -\cos(2\varphi) & 0
\end{pmatrix} \]  \hspace{1cm} (S8)

where \( \varphi \) is the angle of the fast axis. Note that the handedness of CPL in this work is defined from the point of view of the source. Then, the new Stokes vector is:

\[
S_{out} = M_{QWP} \cdot S_{in} = \begin{pmatrix} 1 \\ \cos^2(2\varphi) \\ \sin(2\varphi)\cos(2\varphi) \\ \sin(2\varphi) \end{pmatrix} \]  \hspace{1cm} (S9)

Therefore, the new Stokes parameters are: \( S_0 = 1; \ S_1 = 0.5 + 0.5\cdot\cos(4\varphi); \ S_2 = 0.5\cdot\sin(4\varphi); \ S_3 = \sin(2\varphi). \)
Supplementary note S3

This section analyzes why the infinite $g_{\text{ph}}$ should be only available when the planar nanostructures contain a single mirror symmetry axis. We divide our analysis into two parts: the number of mirror symmetry axes, $N = 0$ and $N \geq 2$.

Case 1: Number of mirror symmetry axes, $N = 0$

With no mirror symmetry axes, the nanostructures are chiral. In this case, none of the elements in Eq. S6 is symmetry-constrained to be vanishing. Therefore, the expression of the photoresponse from a chiral device is in its general form: $V_{\text{ph}} = R_0 \cdot S_0 + R_1 \cdot S_1 + R_2 \cdot S_2 + R_3 \cdot S_3$. Note that the photoresponse discrimination ratio is, $g_{\text{ph}} = |(J_{x,\text{LCP}} - J_{x,\text{RCP}})/(J_{x,\text{LCP}} + J_{x,\text{RCP}})| = |R_3 / R_0|$. Since $R_0$ is usually nonzero when $N = 0$, the $g_{\text{ph}}$ of chiral structures is not symmetry-protected and hence is finite.

Case 2: Number of mirror symmetry axes, $N \geq 2$

When the nanostructure contains two mirror symmetry axes with a relative angle of $\theta$, it is evident that the value of $\theta$ can only be $180^\circ/n$, with $n$ an integer above 2. The nanostructure will also possess an $n$-fold rotation symmetry, $C_n$. Then, the vectorial photoresponse should be invariant under the operation of $C_n$:

$$
\begin{pmatrix}
J_x \\
J_y 
\end{pmatrix} = \begin{pmatrix}
\cos\left(\frac{2\pi}{n}\right) & -\sin\left(\frac{2\pi}{n}\right) \\
\sin\left(\frac{2\pi}{n}\right) & \cos\left(\frac{2\pi}{n}\right)
\end{pmatrix} \begin{pmatrix}
J_x \\
J_y 
\end{pmatrix}
$$

(S10)

The above equation is only possible when both $J_x$ and $J_y$ are zero. Therefore, there is simply no directional photoresponse in the case of multiple mirror symmetry axes.
Supplementary note S4

This section provides a theoretical analysis of the recently discovered near-field chirality (6, 7) via the temporal coupled mode theory and then relates the field asymmetry to the CPL-sensitive directional photoresponse.

S4.1 Analysis of near-field asymmetry using the temporal coupled mode theory

First, we define the two eigenmodes of the plasmonic T-shaped nanostructure as modes A and B (Fig. S1). Mode A possesses a net dipole along the x-axis, while the dipole moment of mode B is along the y-axis. Therefore, mode A and mode B can only be excited by the x- and y-polarized light, respectively.

The mode interference of the plasmonic nanostructure under a specifically polarized illumination can be described by the temporal coupled mode theory 2,3:

\[
\frac{d}{dt} \begin{pmatrix} a \\ b \end{pmatrix} = \begin{pmatrix} j \omega_a - \gamma_{ae} - \gamma_{ai} \\ j \omega_b - \gamma_{be} - \gamma_{bi} \end{pmatrix} \begin{pmatrix} a \\ b \end{pmatrix} + \begin{pmatrix} \kappa_a \\ \kappa_b \end{pmatrix} \begin{pmatrix} E_x \\ E_y \end{pmatrix}
\]  

(S11)

where \(a\) and \(b\) are the oscillation amplitude of the mode A and B. Besides, the mode A and B have respective resonance frequency (\(\omega_a\, \omega_b\)), external decay rates (\(\gamma_{ae}, \gamma_{be}\)), internal decay rates (\(\gamma_{ai}, \gamma_{bi}\), and coupling strength to incident light (\(\kappa_a, \kappa_b\)). \(E_x\) and \(E_y\) denote the complex amplitude of the electric field of the light along the x- and y-axis. Since both

Figure S1. Two eigenmodes of the plasmonic T-shaped nanostructure. +Q and −Q denote oscillating charges in plasmonic mode, with the “+” and “−” signs illustrating the relative phase to incident light.
reflection and transmission are allowed in a single layer nanostructure array, the coupling strength is related to the external decay rates: \( \kappa_a = \sqrt{\gamma_{ae}}; \kappa_b = \sqrt{\gamma_{be}} \).

Considering a circularly polarized illumination, \((E_x, E_y) = E(1, j)\), the mode amplitude can be solved in the frequency domain by replacing \( \text{d/d}t \) with \( j\omega \):

\[
a = \frac{\kappa_a E}{j\omega - j\omega_a + \gamma_{ae} + \gamma_{ai}} \quad \text{(S12)}
\]

\[
b = \frac{j\kappa_b E}{j\omega - j\omega_b + \gamma_{be} + \gamma_{bi}} \quad \text{(S13)}
\]

The local field distribution can then be calculated after knowing the mode profile and amplitude. The electric fields at the left edge, \( E_L \), and right edge, \( E_R \), of the nanostructure are

\[
E_L \propto a - b/2 = -\frac{\kappa_a E}{j\omega - j\omega_a + \gamma_{ae} + \gamma_{ai}} - \frac{j\kappa_b E/2}{j\omega - j\omega_b + \gamma_{be} + \gamma_{bi}} \quad \text{(S14)}
\]

\[
E_R \propto a - b/2 = \frac{\kappa_a E}{j\omega - j\omega_a + \gamma_{ae} + \gamma_{ai}} - \frac{j\kappa_b E/2}{j\omega - j\omega_b + \gamma_{be} + \gamma_{bi}} \quad \text{(S15)}
\]

We then define the field asymmetry, also termed as near-field chirality\(^4\),\(^5\) as:

\[
C = E_R^2 - E_L^2 \quad \text{(S16)}
\]

which is a function of the resonance frequency, decay rates, and the frequency of interest.

The expression of \( C \) can be simplified, because of the relationship \( \kappa_a = \sqrt{\gamma_{ae}} \) and \( \kappa_b = \sqrt{\gamma_{be}} \) as discussed above. In our specific case of mid-infrared plasmonic nanostructures, the internal decay rates of plasmonic nanostructures can usually be regarded as fixed\(^6\). Therefore, the \( C \) is mainly a function of \( \omega, \omega_a, \omega_b, \gamma_{ae}, \gamma_{be} \) as: \( C = C(\omega, \omega_a, \omega_b, \gamma_{ae}, \gamma_{be}) \). Figure S2 shows the maximum available chirality, \( C_{\text{max}} = \max\{C(\omega) : \omega\} \), versus \( (\gamma_{ae}, \gamma_{be}) \) at different normalized frequencies \( \omega_a \) and \( \omega_b \).

Then, we can understand why the near field chirality usually requires an offset in the resonance wavelength of the two modes. As we can see, when \( \omega_b \) is equal to \( \omega_a, C_{\text{max}} \)
vanishes at $\gamma_{ae} = \gamma_{be}$. A small near-field chirality only exists when the $\gamma_{ae}$ and $\gamma_{be}$ are different. In the case of slight off-resonance, $\omega_b - \omega_a = 0.2$, the near-field chirality is the largest. The near-field chirality decreases when the offset further increases, $\omega_b - \omega_a = 0.4$, due to smaller resonance amplitude.

Figure S2. Calculated maximum available near-field chirality at different conditions of $(\omega_a, \omega_b, \gamma_{ae}, \gamma_{be})$. The $C_{\text{max}}$ is defined as $\max\{C(\omega): \omega\}$, where $C$ is the near-field chirality as a function of $(\omega, \omega_a, \omega_b, \gamma_{ae}, \gamma_{be})$. The parameters have been normalized in our temporal coupled mode theory, and $\gamma_{ai} = \gamma_{bi}$ are set to be 0.05. a, At no resonance offset, $\omega_b = \omega_a$, the $C_{\text{max}}$ is small and only available when the two modes possess different decay rates, e.g., a bright mode and a dark mode. b, At slight resonance offset, e.g., $\omega_b - \omega_a = 0.2$, the $C_{\text{max}}$ is the largest, and there are no particular requirements on the decay rates. c, At larger resonance offset, e.g., $\omega_b - \omega_a = 0.4$, the $C_{\text{max}}$ becomes smaller due to the smaller resonance amplitude.

At the optimized condition: $\omega_a = 0.9$, $\omega_b = 1.1$, $\gamma_{ae} = \gamma_{be} = 0.085$, the dependence of $C$ on the normalized frequency is shown in Fig. S3, which further illustrates the origin of near field chirality as the resonance offset between the mode A and B.
Figure S3. Theoretical studies on the near-field chirality induced by mode interference. a, Calculated oscillation amplitude of mode A and mode B. b, Calculated oscillation phase of mode A and B. c, Calculated near field intensity and the resulting asymmetry.

S4.2 Relationship between the near-field asymmetry and directional photoresponse.

The CPL-sensitive vectorial photoresponse will scale with the near-field asymmetry. The reason is the following.

Strictly speaking, the role of the plasmonic nanostructure is twofold. First, it spatially modulates the optical field, leading to an inhomogeneous field profile: $E(x, y)$. Second, the metallic nanostructures spatially modulate the doping level of the graphene sheet $^7, ^8$ and hence the Seebeck coefficients, $S(x, y) ^9$. Upon optical illumination, local photoresponse will be established through the photo-thermoelectric effect, $J(x, y) \propto |E|^2(x, y) \cdot \nabla S(x, y)$. After that, the transport of the local photoresponse is affected by many factors, such as the inhomogeneous conductance of the device channel (see our photovoltage mapping measurement in Fig. S6). As a result, the calculation of the overall photoresponse, $J_{tot}$, is complicated and only possible through numerical modeling. However, we can write the analytic expression of $J_{tot}$ by defining a vectorial local responsivity, $\sigma(x, y)$, as:

$$J_{tot} = \iint_{x,y} |E(x, y)|^2 \cdot \sigma(x, y) dxdy$$  \hspace{1cm} (S17)

Without losing generality, we can assume a rectangular range of integral, $x \in (-x_0, x_0)$ and $y \in (-y_0, y_0)$. Since our plasmonic nanostructure is achiral, the $\sigma(x, y)$ should be parity-
odd regarding the reflection operation, \( x \rightarrow -x \), meaning that \( \sigma(-x, y) = -\sigma(x, y) \). The equation of \( J_{\text{tot}} \) can then be rewritten:

\[
\bar{J}_{\text{tot}} = \iint_{x,y} \left( |\vec{E}(x, y)|^2 - |\vec{E}(-x, y)|^2 \right) \cdot \bar{\sigma}(x, y) dx dy \quad (S18)
\]

In the above equation, the range of integral is halved, \( x \in (0, x_0) \) and \( y \in (-y_0, y_0) \). Note that the term \( |\vec{E}|^2(x, y) - |\vec{E}|^2(-x, y) \) is indeed the near-field asymmetry defined in Eq. S16. Therefore, the CPL-sensitive vectorial photoresponse scales with the near-field asymmetry.
Supplementary note S5

This section presents how we can specifically detect the Stokes parameters, $S_{0,1,2}$. We first write the photoresponse tensor of the up-side-down T-shaped nanostructure (Fig. 2 in main paper):

$$R_1 \sim \begin{pmatrix} 0 & 0 & -0.18 & 0.21 \\ 0.06 & -0.28 & 0 & 0 \end{pmatrix}$$ (S19)

Rotating the nanostructure along with the electrodes by 90°, the new photoresponse tensor would be:

$$R_2 \sim \begin{pmatrix} 0 & 0 & 0.18 & 0.21 \\ 0.06 & 0.28 & 0 & 0 \end{pmatrix}$$ (S20)

The two types of structures can be combined in series. Depending on how we connect the two systems, the photoresponse of the resultant devices can be:

$$\frac{R_1 + R_2}{2} \sim \begin{pmatrix} 0 & 0 & 0 & 0.21 \\ 0.06 & 0 & 0 & 0 \end{pmatrix}$$ (S21)

$$\frac{R_1 - R_2}{2} \sim \begin{pmatrix} 0 & 0 & -0.18 & 0 \\ 0 & -0.28 & 0 & 0 \end{pmatrix}$$ (S22)

Therefore, the $J_x$ and $J_y$ of the first configuration (Eq. S21) can be used for the exclusive detection of $S_3$ and $S_0$, respectively. Besides, the $J_x$ and $J_y$ of the second configuration (Eq. S22) can be used for the specific detection of $S_2$ and $S_1$, respectively.
Figure S4. The thickness of graphene flakes ranges from 2 nm to 3.3 nm, corresponding to few-layer graphene. The text at the left top corner of each image (e.g., C-1A) is the respective label of each flake. The thickness of the metallic structure was measured to be about 75 nm.
Figure S5. Measurement of illumination spot size and estimation of incident power. In our experiments, we used the device with a small footprint (grey box, size < 10 μm) to measure the illumination spot (red area) by spatially scanning the device with a precision stage. We used the Gaussian beam to fit the laser spot on the substrate. The intensity $1/e^2$ radii were measured to be (i) $w_1 = 262 \pm 28 \ \mu\text{m}$ along x-direction and (ii) $w_2 = 215 \pm 7 \ \mu\text{m}$ along y-direction. The black open dots represent the experimental data with a spatial step of 40 μm. Red curves denote the corresponding Gaussian fittings. We always fine-tuned the position of our device to achieve the largest photoresponse, where indicates the center of the laser spot. At this location, the intensity can be calculated as: $I_{(0,0)} = \frac{2P_0}{\pi \cdot w_1 w_2}$, where $P_0$ is the measured overall power from the laser. The incident power received by the device is then: $P_{\text{in}} = I_{(0,0)} \cdot \text{area}$. In particular, the size of our ring device is relatively large, and its incident power is estimated as $P_{\text{in}} = P_0 \left[ \exp\left(-2r_{\text{inner}}^2/w_1 w_2\right) - \exp\left(-2r_{\text{outer}}^2/w_1 w_2\right) \right]$, where $r_{\text{inner}}$ and $r_{\text{outer}}$ represent the inner and outer radius of the ring-shaped device, respectively.
Figure S6. Studies on the directional nonlocal photoresponse in graphene via photovoltage mapping. a, Optical microscopy image of the device with metallic structures sitting on the graphene flake. To achieve good spatial resolution, we have intentionally used enlarged metallic structures (as five times large as the structure in the main paper) and a visible laser (532 nm wavelength) with a focus spot size of about 300 nm. At these conditions, plasmonic effects do not play a role here, so that we can probe the local responsivity. b, In-situ mapping of Raman intensity measured at the 2D band of graphene (Raman shift 2685 cm⁻¹). Red regions denote the graphene area, while the blue regions are covered with metal or SiO₂ substrate. c, In-situ mapping of photovoltage measured at zero external bias. Red and blue colors represent positive and negative open-circuit voltages. Notably, the T-shaped metallic structures also contribute large photoresponse, although they are not connected to the device electrodes. The directional and nonlocal character of the photoresponse in graphene can be captured by the Shockley-Ramo theorem¹⁰. d, Photovoltage mapping around the vertical T-shaped structure with a higher resolution. The incident power is 15 mW. 100x objective was used. Our results show that the responsivities at the left and right ends of the T-shaped metallic structure are opposite. Therefore, the broken symmetry in the near field of nanostructure will lead to an unbalanced photoresponse.
Figure S7. Exclusion of graphene plasmon effects in our device. a-d, Simulated near field profile at different Fermi levels of graphene, with the incident light at 4 µm wavelength. We have used the surface conductivity material model to incorporate graphene into our simulation (FDTD, Lumerical). We set the conductivity scaling number to 3, corresponding to trilayer graphene. The scattering rate in our simulation is the default value of 0.00051423 eV. As we can see, graphene plasmons are only efficiently generated when the Fermi level is above 0.3 eV. e, Estimated Fermi level in graphene as a function of applied gate voltage. $V_{\text{dirac}}$ denotes the gate voltage when the Fermi level crosses the Dirac point, namely, the charge neutral point. In our experiments, the $V_{\text{dirac}}$ is at 136 V (Fig. S20). Therefore, the Fermi levels of the graphene in our device span from -0.21 eV ($V_g = 0$V) to 0.12 eV ($V_g = 180$V). Therefore, the Fermi levels in our experiments (wavelength at 4 µm) are still too low to efficiently excite graphene plasmons.
Figure S8. Simulation process to calculate the near-field induced photoresponse. We start from (a) the layout of a nanostructure unit. b, Because of the doping effect of metals on graphene, there exists a Seebeck coefficient gradient \( \nabla S(r) \) at the metal-graphene interfaces, where \( r \) denotes the position. c, Simulated near-field intensity of plasmonic nanostructures under RCP illumination. Note that the local absorption of graphene scales with the field intensity; we assume that the inhomogeneous heating of graphene electrons is similar to the field profile, \( \Delta T(r) \propto |E(r)|^2 \). This assumption can also be revised in the future to capture the saturation effect. d, From the Seebeck coefficient gradient and electron temperature levitation profile, we can derive the inhomogeneous driving forces in the structure, \( F(r) \propto \nabla S(r) \cdot \Delta T(r) \). e, The current flow due to the local photoresponse is calculated by solving the Navier-Stokes equations where the electrons in graphene are regarded as fluid. \( \rho \) represents the effective pressure of the fluid in our simulation.
Figure S9. Experimental studies on the geometry of metallic nanostructures for CPL sensitivity. a, SEM image of four devices, in which the metallic nanostructures have a constant $h$ of 1100 nm and different $v$ of 500 nm, 600 nm, 750 nm, and 900 nm. The four devices were fabricated on the same patterned graphene sheet for a fair comparison. (b-e) measured QWP angle-dependent photovoltages in the devices with (b) $v = 500$ nm, (c) $v = 600$ nm, (d) $v = 750$ nm, and (e) $v = 900$ nm. f, Simulated CPL sensitivities for various $h$ and $v$, a replica of Fig. 2b in the main paper. The geometric parameters of the four fabricated devices are denoted with four stars. g, Extracted experimental CPL sensitivities of the four devices, confirming our numerical calculation. Therefore, we choose the optimized design at $v = 750$ nm and $h = 1100$ nm.
Figure S10. Simulated near field distribution under LCP illumination, showing that the field intensity at the left edge is about six times larger as the right edge. The field intensity along the white dashed line in (a) is plotted in (b). The normalized peak intensities at the left and right edges of the nanostructure are 725 and 118, respectively.
Figure S11. Experimental studies on the CPL-insensitive Jy in the mirror-symmetric nanostructure. a, SEM image of the fabricated device, where electrodes were placed in such a way to characterize the Jy as defined in Fig. 2a in the main paper. b, Measured QWP angle-dependent photovoltages and the attributed contributions from the four Stokes parameters, S0,1,2,3. Inset shows the respective responsivities, R0,1,2,3. Note that we simulated the short-circuit currents, Isc, of our device but measured the open-circuit voltages, Voc, in our experiments. The Isc and Voc are linearly related but with opposite signs in our device.
Figure S12. Simulation of the polarization-dependent photocurrents to extract the photoresponse tensor. The polarization states of the incident light are (a) $S_{0,1,2,3} = (1, 1, 0, 0)$; (b) $S_{0,1,2,3} = (1, -1, 0, 0)$; (c) $S_{0,1,2,3} = (1, 0, 1, 0)$; (d) $S_{0,1,2,3} = (1, 0, -1, 0)$; (e) $S_{0,1,2,3} = (1, 0, 0, -1)$; (f) $S_{0,1,2,3} = (1, 0, 0, 1)$. Averaged $J_x$ and $J_y$ denote the normalized photocurrents flowing along $x$ and $y$ directions, respectively. Therefore, the normalized photoresponse tensor can be obtained by solving the following relation.

$$
\begin{pmatrix}
0 & 0 & -0.18 & 0.18 & -0.21 & 0.21 \\
-0.22 & 0.34 & 0.06 & 0.06 & 0.08 & 0.08
\end{pmatrix} \sim \begin{pmatrix} R_{x0} & R_{x1} & R_{x2} & R_{x3} \\
R_{y0} & R_{y1} & R_{y2} & R_{y3}\end{pmatrix} \begin{pmatrix} 1 & 1 & 1 & 1 & 1 & 1 \\
-1 & -1 & -1 & -1 & -1 & -1
\end{pmatrix}
$$

(S23)

And we get:

$$
\begin{pmatrix} R_{x0} & R_{x1} & R_{x2} & R_{x3} \\
R_{y0} & R_{y1} & R_{y2} & R_{y3}\end{pmatrix} \sim \begin{pmatrix} 0 & 0 & -0.18 & 0.21 \\
0.06 & -0.28 & 0 & 0 \end{pmatrix}
$$

(S24)
Figure S13. Measured polarization-dependent photovoltages at different orientation angles of the rectangle device with an interval of 15°.
Figure S14. Measured polarization-dependent photovoltages of the segments of a curved device. (a) Optical microscopy and (b) SEM images of the fabricated curved device, which consists of 6 segments named D1, D2, D3, D4, D5, and D6. Electrodes were fabricated to probe the photoresponse of each segment. c, Zoom-in images of the six segments. Insets show an amplified view of the nanostructures. d, Measured photovoltages of the six segments and their summation, reinforcing our results in Figure 3 in the main paper.
Figure S15. Half-ring cascaded device. a, SEM image of the fabricated device, where the orientation of T-shaped nanostructures, $\theta$, spans from 0° to 180°. b, Measured QWP angle-dependent photovoltages and the attributed contributions from the four Stokes parameters, $S_{0,1,2,3}$. Inset shows the respective responsivities, $R_{0,1,2,3}$. 
Figure S16. Repeatable enhancement of photoresponse in graphene ribbons device. Other than the data presented in the main paper, here we provide another two pairs of devices to fairly compare the graphene sheet and ribbons devices. a,b, Optical microscopy image and measured photoresponse of the second pair of devices. The incident power is about 69.2 μW. c,d, Optical microscopy image and measured photoresponse of the third pair of devices. The incident power is about 793.9 μW. Note that this device is not passivated with 10 nm Al₂O₃ layer by ALD, which usually shows lower responsivity based on our experience. The enhancement factors are about 5.
Figure S17. Measurement of the noise spectra in the dark condition. The noise of our device was measured with a lock-in amplifier (see methods for details). The Johnson noise is calculated as \( \sqrt{4k_B TR\Delta f} \), where \( k_B \) is Boltzmann’s constant in joules per kelvin, \( T \) is the device’s absolute temperature in kelvins, \( R \) is the resistance of the device, and \( \Delta f \) is the bandwidth in hertz over which the noise is measured. The resistance of our ribbons device is 2095 \( \Omega \) at no gate voltage, so that the resulting Johnson noise limit is about 6 nV Hz\(^{-1/2}\). In our experiments, the noise density reaches a saturated value of about 10 nV Hz\(^{-1/2}\) when the frequency is above 100 Hz.
Figure S18. Studies on the power-dependent noise and detectivity. The measured device noise becomes linearly dependent on the incident power above 1 μW, which has not been fully understood by the authors yet. Note that the detectivity of ellipticity is $D_{\text{ellipticity}} = \frac{(180^\circ/\pi) \cdot N_{\text{light}}}{(2 \cdot R_3 \cdot P)}$, thereby the linearly increasing noise, $N_{\text{light}}$, will compensate the linearly increasing signals, $R_3 P$, lead to a saturated detectivity at about $0.03^\circ \text{Hz}^{-1/2}$. 
**Figure S19. Characterization of frequency response.** Based on the response time of about 900 ns (as shown in Fig. S20 below), the 3dB frequency of our device, $f_{3dB}$, can be estimated as $0.35/900$ ns $\approx 390$ kHz. We measured almost constant photoresponse up to 4 kHz in our experiments.

$$V_{ph} = \frac{V_0}{(f/f_{3dB})^2 + 1}$$
Figure S20. **Time response characterized with a pulsed laser.** The orange shade indicates the pulse illumination with a duration of 1000 ns. The rise time of the pulsed laser is below 10 ns. The fitted rise time, $\tau_{\text{rise}}$, and fall time, $\tau_{\text{fall}}$, are 886 ns and 902 ns, respectively.
Figure S21. Dependence of response time on the device geometry. a-c, The layout of three device with different device lengths, $L$, and widths, $W$. Graphene sheets are not shown in the layout. Insets show the optical microscopy images of the fabricated devices. d-f, Measured time responses for the three devices. The orange shade indicates the pulse illumination with a duration of 1000 ns. The extracted fall times scales with the aspect ratio of our device, $\square = L/W$. Note that the cascading of nanostructures will increase the resistance and decrease the capacitance of our device proportionally, so that the RC delay should be constant if the response is limited by the intrinsic photoresponse speed of nanostructures. However, this is not the case. The measured response time will scale with the resistance of our device and hence the aspect ratio, only when the overall capacitance is independent of the device geometry, for example, dominated by parasitic capacitance in the setup. Therefore, we speculate that the measured response time is setup-limited. The intrinsic photoresponse speed should reach the sub-ns level due to the fast relaxation of hot electrons in graphene$^{13,14}$. 
Figure S22. Gate-dependent resistance and extracted mobility of our ribbons device.
At a constant drain-source voltage, $V_d = 1 \text{ mV}$, we measured a peak resistance at the gate voltage of 136 V, where the Fermi level of graphene reaches its charge neutral point. We extract the hole mobility from the orange region to be 2446 cm$^2$V$^{-1}$s$^{-1}$. 
Figure S23. Measured photoresponse versus incident power. The power-dependent photoresponse can be fitted with a power function: $V_{oc} = A \cdot P^n$, where $A$ and $n$ are fitting parameters. $V_{oc}$ in units of mV, and $P$ in units of mW. The illumination is RCP light. The fitted $n$ is below 1, probably due to the screening effect observed in previous works\textsuperscript{14}. 

$V_{oc} = 63.3 \cdot P^{0.82}$

$V_{oc} = -6.1 \cdot P^{0.84}$
Figure S24. Calculated wavelength dependence of our CPL photodetector. The photoresponse is optimized at 4 μm wavelength. The full width at half maximum (FWHM) is about 1 μm. Notably, our device is exclusively responsive to CPL at all wavelengths because of its symmetry.
Figure S25. Calculated fill factors of (a) the ring-shaped device, (b) proposed meander ring-shaped device, and (c, d) proposed cascaded rectangle devices. The fill factors of the ring-shaped and meander-shaped devices are the division of device area by the area of the rectangle with dashed borders.
Figure S26. Near-field chirality and super-chiral field. Comparison of the (a, c) near-field chirality defined in this work and (b, d) the super-chiral field under the (a, b) RCP and (c, d) LCP illumination. The super-chiral field has been normalized by the chirality of the incident circularly polarized light, and the enhanced chirality as $C = \frac{\text{Im}(E^* \cdot H)}{\text{Im}(E_0^* \cdot H_0)}$, where $E_0$ and $H_0$ are the electric and magnetic fields of the incident light.
Figure S27. Simulated absorption spectra of metallic nanostructures ($A_{Metal}$) and graphene ($A_{Graphene}$). The graphene layer number is 5, and the relative chemical potential is 0 eV. The absorption of graphene is about 30%.
| No. | Description                  | Wavelength Range (µm) | Responsivity (V/W) | Bias Voltage (V) | Ref. (with hyperlink) | Remarks                                      |
|-----|------------------------------|-----------------------|--------------------|-----------------|----------------------|----------------------------------------------|
| 1   | Germanium photodiodes        | 0.8 – 1.8 (NIR)       | 3400               | 3 V             | FDG50, Thorlabs       | 0.85 A/W; Shunt Resistance=4000Ω             |
| 2   | Germanium photodiodes        | 0.8 – 1.8 (NIR)       | 2160               | 0.25 V          | 460004-2, Teledyne    | 0.9 A/V; Resistance=2400Ω                   |
| 3   | Germanium photodiode         | 1.5 (NIR)             | 10                 | 0-2 V           | Nat. Photon. 15, 925 (2021) | I_{sc}=0.7mA@4mW; Resistance=59Ω             |
| 4   | MCT photovoltaic detector    | 2.7 – 5 (MIR)         | 100                | 0-0.8 V         | VL5T0, Thorlabs       | 1 A/W; Resistance~100Ω                     |
| 5   | InAsSb photovoltaic detector | 4 – 5.9 (MIR)         | 21                 | 0 V             | P11120-201, Hamamatsu | 1.6 A/W; Shunt Resistance~13Ω               |
| 6   | Thermopile detector          | 0.19 – 20 (MIR)       | 0.1                | 0 V             | TD10X, Thorlabs       |                                             |
| 7   | Thermopile detector          | 3 – 5 (MIR)           | 50                 | 0 V             | T11361-01, Hamamatsu  |                                             |
| 8   | Nanostructure/Graphene detector | 4 (MIR)              | 392                | 0 V             | This work             | Functional and CMOS-compatible              |
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