We report experimental studies of bright-state polaritons of four-wave mixing (FWM) and six-wave mixing (SWM) signals through cascade nonlinear optical parametric amplification processes in an atom-cavity composite system for the first time. Also, the coexisting cavity transmission modes of parametrically amplified FWM and SWM signals are observed. Finally, electromagnetically induced absorption by the FWM cavity modes in the probe beam is investigated. The investigations can find potential applications in multi-channel narrow-band long-distance quantum communication.

Cavity quantum electrodynamics has been intensively studied in recent years, owing to the important applications in quantum optics and nonlinear optics. When a two-level atom is placed inside the optical cavity under the strong coupling condition, the “normal-mode splitting” have been experimentally demonstrated in many atomic systems. If the intracavity medium is a coherently-prepared three-level atomic medium, the so-called “dark-state polariton” peak appears in the cavity transmission spectrum (CTS) due to the intracavity electromagnetically induced transparency (EIT). Such “dark-state polariton” have potential application in the long-lived storage of quantum information and quantum computation. Moreover, when the atomic system is in the double-A configuration, correlated photon pairs with high generation rate and narrow bandwidth have been generated in the cold atomic ensemble by using an optical parametric amplification (OPA) process operated below its oscillation threshold. Meanwhile, bright correlated anti-Stokes and Stokes twin beams with hot atoms inside an optical cavity have also been obtained by an OPA process above the threshold. When an ensemble of cold atoms strongly coupled to a high-finesse optical cavity, and the control light can be generated from the vacuum, then the vacuum-induced transparency was observed. In addition, four-wave mixing (FWM) and six-wave mixing (SWM) processes based on third-order and fifth-order nonlinear processes in EIT media have also attracted lots of attention in recent years, in which a strong coupling beam renders a resonant, opaque medium nearly transparency while enhancing the nonlinearity, and the coexistence of these two nonlinear processes due to double EIT windows and atomic coherences has been reported. The nonlinear process plays important role in entanglement generation and cascade-nonlinear optical process. To the best of our knowledge, the OPA seeded with FWM or SWM signal in an atom-cavity composite system has not been reported, and the investigation of it would be important for building multi-channel nonlinear optical devices and ultra-narrow linewidth photon sources for long-distance quantum communication.

In this letter, we report our investigations of bright-state polaritons of FWM and SWM signals through cascade OPA process in an atom-cavity composite system. The bright-state polaritons of FWM signal with 5 MHz linewidth are obtained. We also report the coexisting cavity modes of parametrically amplified FWM and SWM signals in an inverted-Y-type atomic system for the first time. Moreover, the electromagnetically induced absorption (EIA) peaks induced by the multiple cavity modes of the FWM signal are observed in the probe beam. These results are well explained by the presented theoretical model. The investigation will help us to better understand the interactions between the strongly coupled multi-level atoms and the optical cavity, which can find application in quantum information processing.
Results

Fig. 1(a) shows the experimental setup, while the energy levels of the atomic system used in the experiment is shown in Fig. 1(b1), where the fields $E_1$ (frequency $\omega_1$, wave vector $k_1$, Rabi frequency $G_1$, and wavelength 780.2 nm), $E_2$ ($\omega_2$, $k_2, G_2$, and 776.16 nm) & $E_3$ ($\omega_3$, $k_3, G_3$, and 780.2 nm) & $E_1$ ($\omega_3$, $k_1, G_1$) are used as probe field, pumping fields and coupling fields, respectively. The resonant transition frequencies of $|0\rangle \rightarrow |1\rangle$, $|1\rangle \rightarrow |2\rangle$ and $|1\rangle \rightarrow |3\rangle$ are $\Omega_1$, $\Omega_2$ and $\Omega_3$, respectively. Then the frequency detuning for each field can be defined as $\Delta_i = \Omega_i - \omega_i$ (i = 1, 2, 3). With all three beams present in Fig. 1(b1), three phase-conjugate signals (FWM signals $E_{FSR}$ & $E_2$ and SWM signal $E_3$), satisfying the phase-matching conditions of $k_{F1} = k_1 + k_f - k_2$, $k_{F2} = k_1 + k_3 - k_2$, and $k_3 = k_1 + k_2 - k_1 - k_3$, can be generated simultaneously at the center of the atomic cell and propagate along the optical axis of the cavity (dashed line in Fig. 1(a)). Here, the narrow signals $E_{FSR}$ and $E_2$ are generated within the EIA window of $\delta = 0$, where $\delta = \Delta_1 + \Delta_2$ is two-photon detuning, while the broad signal $E_3$ will be generated due to no assistance of Doppler-free atomic coherence and EIA window. However, if the $E_1$ beam has a sufficiently high power, as well as is far detuned from $|0\rangle \rightarrow |1\rangle$, a spontaneous parametric FWM (SP-FWM) process will occur in the degenerate two-level atomic configuration (Fig. 1(b2)), which generates two weak fields (Stokes field $E_5$ and anti-Stokes field $E_{ASS}$), satisfying $2k_1 = k_{SS} + k_{ASS}$ (Fig. 1(b3)), on a forward cone. The generated $E_{FSR}$ (or $E_2$) signal is naturally injected into the input Stokes port of the SP-FWM process, and is parametrically amplified, where the process will serve as an OPA. The parametrically amplified signal denoted as $E_{F1}^{SS}$ (or $E_2$) is still generated at the center of the atomic cell and propagate along the optical axis of the cavity, so they are mode-matched to the cavity and form the cavity modes.

According to the expression of $a_{F1}^{SS}$ given in the Method part, Fig. 1(c) shows the calculated normalized CTS versus $\Delta_{ac}$ (defined as $\Delta_{ac} = \Omega_1 - \omega_{ac}$, with $\omega_{ac}$ being the resonant frequency of the cavity) and $\Delta_1$, which exhibits a double-peak structure along the $\Delta_1$ direction. This structure indicates the cavity polaritons of the parametrically amplified field $E_{F1}^{SS}$, which results from the coupling between the cavity mode of $E_{F1}^{SS}$ and atoms with strength $g_N\sqrt{N}$, where $g$ is the single-atom-cavity coupling strength, and $N$ is the number of atoms in the cavity. When $\Delta_1$ changes from negative to positive continuously, the polaritons move, that is dictated by the condition of $\Delta_1 - \Delta_{ac} = 0$, and finally leads to the avoided-crossing plot in Fig. 1(c). Fig. 1(d) is plotted with the same variables as Fig. 1(c) but in a wider range, where multiple polaritons at a fixed $\Delta_1$ can be seen, and their positions change as $\Delta_1$ changes. It is found that $\Delta_{FSR} - \Delta_{ac} = l \omega_{FSR}$ (l is an integer, and $\omega_{FSR}$ is the free-spectral range (FSR) with medium) is always fulfilled at all the polaritons, i.e. $\Delta_1 - \Delta_{ac} = l \omega_{FSR}$ (named cavity transmission window) is satisfied for all the polaritons. So, when $\Delta_1$ changes, the corresponding $\Delta_{ac}$ for each transmission polaritons of $E_{F1}^{SS}$ changes accordingly. We also show normalized CTS versus $\Delta_2$ and $\Delta_{ac}$ in Fig. 1(e), in which the polaritons nearly do not move with $\Delta_2$. The reason is that the energy shifts induced by dressing field $E_2$ at different $\Delta_1$ values make the window $\Delta_1 - \Delta_{ac} = l \omega_{FSR}$ not move.

It is worth mentioning that the cavity polaritons of the parametrically amplified field $E_{F1}^{SS}$ and $E_2$ will have ultra-narrow linewidths, which would be useful for long-distance quantum communications. On the one hand, the linewidths of $E_{F1}^{SS}$ and $E_2$ are narrowed by the EIA window, on the other hand, $E_{F1}^{SS}$ and $E_2$ form cavity modes, and then the cavity polaritons will be further narrowed due to the large dispersion change and reduced absorption accompanying EIA.

Parametrically amplified bright-state polarization of FWM signal

We first measure the CTS without $E_2$ and $E_3$. In this case, the SP-FWM process generates a Stokes field $E_5$ (measured and shown in Fig. 2(a1)) and an anti-Stokes field $E_{ASS}$ (measured and shown in Fig. 2(a2)) on a forward cone. Such a process can also act as an OPA for the $E_5$ in (Fig. 2(a3)) injected into the Stokes port (in Fig. 2(a4)). The parametrically amplified signal denoted as $E_{F1}^{SS}$ forms cavity mode, which couples with the atoms. By scanning $\Delta_1$ across the transition $F = 3 \rightarrow F'$ in $^{87}$Rb at different $\Delta_{ac}$ values and taking $\Delta_2 \approx 1.2$ GHz, the measured CTS of $E_{F1}^{SS}$ versus $\delta$ are shown in Figs. 2(b1)–(b3), where the lower curves (ii) and top curves (i) are the CTS of $E_{F1}^{SS}$ and the corresponding EIA window, respectively, with the atomic cell temperature $T \approx 77$ C, and powers of $E_1$, $E_2$, and $E_3$, $P_1 = 8$ mW, $P_2 = 22$ mW, $P_3 = 22$ mW, respectively. The CTS of $E_{F1}^{SS}$ without the atom-cavity mode coupling is shown in Fig. 2(b2) by the curve (iii), where one dip and two peaks can be seen clearly. The two peaks represent the Atuler-Townes splitting of the signal $E_{F1}^{SS}$, which derive from the two dressed states, namely, bright states. The dip comes from the dark state induced by the destructive interference between two dressed states. Tuning $\Delta_{ac}$
with the atom-cavity mode coupling indicates that, when the $\Delta_1 - \Delta_{ac} = 0$ window is overlapped with the EIA window, the two peaks reach their maxima simultaneously on the curve (ii) of Fig. 2(b2). Comparing the curve (iii) with the curve (ii), the dip (corresponding to dark state) and the two peaks is amplified, which stems from the dressing splitting of the atom-cavity mode coupling to the energy level $|1\rangle$. So the two peaks on the curve (ii) represent the cavity polaritons of $E_{P1}$, as shown in Fig. 1(c). When the $\Delta_1 - \Delta_{ac} = 0$ window deviates from EIA window by decreasing $\Delta_{ac}$, the two peaks become asymmetric, and the left peak is amplified by the atom-cavity mode coupling as shown in Fig. 2(b1). The left peak corresponds to one of the bright states, so one intracavity bright state (“bright-state polariton”) with a linewidth about 12 MHz is obtained. The other narrow “bright-state polariton” (~5 MHz) is obtained when increasing $\Delta_{ac}$ in Fig. 2(b3). The calculated CTS curves of $E_{P1}^s$ ($\propto |\psi_i\rangle^2$), which agree well with the measured results, are shown in Figs. 2(b4)–2(b6).

By scanning the voltage imposed on PZT, Fig. 2(c1) displays the parametrically amplified bright-state polaritons of $E_{P1}^s$ versus $\Delta_{ac}$ at different $\delta$ values, which is set as $-12$ MHz, $0$ and $12$ MHz from bottom to top by taking different $\Delta_1$ with fixed $\Delta_2$. Each CTS curve exhibits two bright polaritons, satisfying $\Delta_1 - \Delta_{ac} = \text{fsweep}$, having a separation of $\omega_{P1}^s$, and changing their positions with $\Delta_1$, as predicted by the theoretical result in Fig. 1(d). Also, the heights of the polaritons change with $\Delta_1$, since the polaritons are enhanced at two-photon resonant. Next, by fixing $\Delta_1$ and taking different $\Delta_2$, the bright-state polaritons of $E_{P1}^s$, versus $\Delta_{ac}$ with $\delta$ set as $-15$ MHz, 0 and 15 MHz from bottom to top are shown in Fig. 2(c2), where the heights of the polaritons change with $\Delta_2$, due to two-photon detuning, but the positions are fixed. That is similar to the case shown in Fig. 1(e).

The polaritons of $E_{P1}^s$ can be affected by the temperature of atomic cell, since increasing $T$ can increase the atomic number $N$, and then enlarge the collective coupling factor $g\sqrt{N}$, which will yield dressing to the polaritons. According to the theoretical model presented in the method part, the splitting eigenvalues induced by the dressing of atom-cavity mode coupling to the polaritons can be given by $\lambda_{1,2} = (-\Delta_1 \pm \sqrt{\Delta_1^2 + 4g^2N})/2$, which correspond to the positions of the two polariton peaks in the CTS curves. So the splitting of the two polariton peaks will increase with $T$. Also, increasing temperature will lead to increased atomic number participating in the phase-conjugate FWM process as well as OPA process, which will enhance the height of the two polariton peaks. The measured CTS curves versus $\delta$ with increasing $T$ from bottom to top are shown in Fig. 3(a1), where not only the height but also the splitting of the two polariton peaks increase with $T$. Fig. 3(A1) shows the calculated results corresponding to Fig. 3(a1), according to the expression of $a_{P1}$. The experimentally measured (squares) and theoretically calculated (solid line) splittings of polaritons versus $T$ are shown in Fig. 3(a2), where the splitting indeed increases with $T$. The experimental results agree with the theory.

The polaritons of $E_{P1}^s$ can also be influenced by the powers $P_1$ and $P_2$. The measured CTS curves versus $\delta$ with increasing $P_2$ from bottom to top are presented in Fig. 3(b1). The height and the splitting of the two polariton peaks become more pronounced as $P_2$ increase. Figure 3(b2) shows the splitting versus $P_2$ with the atom-cavity mode coupling (upper plots) and without the atom-cavity mode coupling (lower plots), where squares and dots are experimental results, while the solid curves are the corresponding fittings. The splitting on upper plots in Fig. 3(b2) is mainly induced by the dressing of the atom-cavity mode coupling and field $E_2$ to polaritons, while the splitting on lower plots mainly results from the dressing of the field $E_2$. By comparing the two cases, the dressing splitting of the atom-cavity mode coupling can be obtained. While the dressing splitting induced by field $E_2$ to the polariton can be seen by increasing $P_2$. Similar investigation has been done for the influence of $P_1$, in Figs. 3(c1) and 3(c2). However, the upper squares (or solid line) in Fig. 3(c2) is much easier to saturate than that in Fig. 3(b2). The dressing splitting of the polariton in Fig. 3(c2) is mainly determined by the atom-cavity coupling and field $E_2$ when the power of $P_1$ is weak, which is similar to the case in Fig. 3(b2). As $P_1$ is increased, the dressing of field $E_2$ gets larger and has to be considered. So the dressing splitting of field $E_2$ to polaritons is included in Fig. 3(c2) compared with Fig. 3(b2). Therefore, the saturation behavior with increasing $P_1$ in Fig. 3(c2) mainly results from the balanced interactions between the destructive and constructive interferences of different dressing pathways induced by the atom-cavity mode coupling, field $E_1$ and field $E_2$.

### Coexisting cavity modes of parametrically amplified FWM and SWM signals

If all laser beams are opened except $E_2$ (Fig. 1(b1)), we can get an inverted Y-type system, in which there will be $E_S$ and $E_P$ but no $E_{F1}$. As indicated by Fig. 4(f), $E_S$ signal is naturally used as a seed injected into the Stokes port of SP-FWM process and amplified as $E_{F1}^s$. Then, the $E_{F1}^s$ signal, mode-matched to the cavity, forms a cavity mode. In free space, the measured coexisting spectrum of $E_{P1}^s$ and $E_{F2}$ versus $\delta$ with $\Delta_2 = 1.2$ GHz is displayed as the bottom curve in the inset of Fig. 4(a), in which a narrow peak labeled as “S” (corresponding to EIA dip labeled as “E” on the top curve) and a broad signal can be seen. The frequency of peak “S” changes with $\Delta_2$, so peak “S” is the signal $E_{F2}$, while the broad signal is $E_{F2}$ without the EIA window. With both $E_{P1}^s$ and $E_{F2}$ resonant in the cavity, the measured CTS versus $\delta$ is given by the bottom solid curve in Fig. 4(a), where we can see five peaks labeled as “P1”, “P2”, “ST”
“P3”, and “P4”, respectively, and the peak “ST” corresponds to the EIA dip “E” on the top dash-dotted curve. Without $E_2$, both the “ST” peak and “E” dip disappear in Fig. 4(b). So we can identify the “ST” peak is the transmission peak of $E_2^d$, and the other four peaks are the cavity modes of $E_2^d$ picked out by the cavity windows from the broad gain.

Tuning the cavity length to make a cavity transmission peak of $E_2^d$ overlap with the “ST” peak (i.e., the cavity window overlaps with the EIA window), the resulted CTS (the solid curve) and the corresponding EIA (the dash-dotted curve) versus $\delta$ are shown in Fig. 4(c), where the coexisting cavity modes of the SWM ($E_2^d$) and FWM ($E_2^d$) signals are obtained. By blocking $E_2$ at this time, the cavity mode of the SWM signal disappears, and only the cavity modes of $E_2^d$ survive (dashed line of Fig. 4(c)). Comparing the results with and without $E_2$ in Fig. 4(c), the enhancement (bright-state polariton of $E_2^d$) due to the coupling of cavity mode of $E_2^d$ with atoms can be seen. Also, the pulling of the cavity resonant frequency due to the dispersion change induced by $E_2^{20.21}$ can be observed. By fixing $\Lambda_1$ and $\Delta_2$ at the point of coexisting cavity modes and scanning $\Delta_{ac}$ in Fig. 4(d), the measured coexisting cavity modes of $E_2^d$ and $E_2^d$ versus $\Delta_{ac}$ are shown on the solid curve, while the dashed curve shows the cavity modes of $E_2^d$ without $E_2$. Also, the pulling of the cavity resonant frequency is observed in Fig. 4(d). According to the coupled atom–cavity model, the CTS amplitude of $E_2^d$ is given by $\alpha_{2c} = T_{FS}(\alpha_{2c}^{ph})/\{[(\Delta_2 - \Delta_{ac}) + i\gamma] + \omega^2N/d_1 + |G_2|^2/d_2 + |G_3|^2/d_3\}$, where $\Delta_2 = \Omega_1 + i(\Delta_1 - \Delta_2)$, and $G_2^2$ is the Rabi frequency of $E_2^d$. The transmission coefficient of intensity of $E_2^d$ is $T_{FS} = (t_2t_3)^2e^{-2\alpha_2}/\{[(1-r_2r_3e^{-2\alpha_2})^2 + 4r_2r_3e^{-2\alpha_2}\sin^2(\phi/2)]\}$, where $r_i$ ($t_i$) is the reflection (transmission) coefficient of the mirror $M_i$ of the cavity with $r_1^2 + t_1^2 = 1$, and $\phi(\omega_{2c}) = 2\pi(\Delta_{ac} - \Delta_2)/\omega_{FSR} + (n - 1)\omega_{FSR}/c$ is the round-trip phase shift experienced by $E_2^d$, with light speed in vacuum $c$, length of the atomic cell $L_{ac}$, and FSR of the empty optical cavity $\omega_{FSR} = c/L_{ac}$ with a cavity length $L_c$. The terms $\gamma = 2(\omega_{2c}/c)\Im[(1 + \chi)^{1/2}]$ and $n = \Re[(1 + \chi)^{1/2}]$ are the intensity absorption coefficient and refractive index of the atomic medium, respectively, with $\chi = 2\omega_{2c}^2N_{ac}/[\omega_{FSR}(d_1 + |G_2|^2/d_2 + |G_3|^2/d_3)]$ when only the linear susceptibility is considered. Using the expressions of $a_2$ and $T_{FS}$, the calculated normalized CTS of coexisting $E_2^d$ and $E_2^d$ versus $\delta$ are shown on the bottom curve in Fig. 4(e), which agrees with the experimental result.

Besides, the EIA peaks induced by the multiple cavity modes of $E_2^d$ are observed in Fig. 4(b), where each inset is a measured EIA peak multiplied by 10. Because each cavity mode of $E_2^d$ couples to $|0\rangle\rightarrow|1\rangle$ and propagates along the same direction as $E_1$, the two-photon resonance (one photon from $E_1$ and another from $E_2^d$) induces the Doppler-free atomic coherence$^{23}$. Owing to the far detuning of the cavity modes of $E_2^d$, the EIA is induced instead of EIT. The probe absorption can be obtained by $\rho_1^{(1)} = iG_1/\{[d_1 + |G_2|^2/d_2 + |G_3|^2/d_3]/d_1 + |G_2|^2/d_2 + |G_3|^2/d_3\}$ with $d_1 = i(\Delta_1 - \Delta_{ac}) + i\Gamma_1$, $d_2 = i(\Delta_2 - \Delta_{ac}) + \Gamma_1$, and $d_3 = i(\Delta_3 - \Delta_{ac}) + \Gamma_1$, which indicates $E_2^d$ field and the cavity modes of $E_2^d$ are not easy to be observed separately, because they overlap with the EIA created by the $E_2^d$ field, which is much stronger than the cavity modes. The induced EIA (insets of Fig. 4(b)) by the cavity modes of $E_2^d$ can be observed by tuning $\Delta_{ac}$ or $\Delta_2$ to separate it from the EIA of $E_2^d$. The calculated probe absorption by $\rho_1^{(1)}$ is shown on the top curve in Fig. 4(e), where the highest EIA peak is the sum from
E$_2$ and cavity modes of E$_{2}^{\text{CTS}}$ & E$_{3}^{\text{CTS}}$, and the other peaks are induced by the cavity modes of E$_{2}$.

**Discussion**

We have experimentally studied the bright-state polarizations of multi-wave mixing signals through OPA processes in a multi-level atom-cavity composite system. It is demonstrated that the polarizations are much narrowed due to EIA and cavity window. It would be important for the narrow-band long-distance quantum communications. Besides, coexisting cavity modes of the generated SWM and FWM signals are also observed, which can help us to better understand the interactions between the strongly coupled multi-level atoms and the optical cavity. Moreover, the induced EIA by the multiple FWM cavity modes in the probe spectrum is observed. Such investigation in an atom–cavity coupling system may find potential applications in building multi-channel nonlinear optical devices for quantum information processing.

**Methods**

**Experiment setup.** The atom-cavity composite system is shown in Fig. 1(a), where a 7 cm long Rb atomic cell is placed inside a 38 cm long optical ring cavity. The concave mirrors M1 and M2 (with the same radius of curvature of 100 mm) have 99.9% and 7 cm long Rb atomic cell is placed inside a 38 cm long optical ring cavity. The concave mirrors M1 and M2 (with the same radius of curvature of 100 mm) have 99.9% and 7 cm long Rb atomic cell is placed inside a 38 cm long optical ring cavity. The concave

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**Acknowledgments**

This work was supported by the NRBFC (2012CB921804), NSFC (61205112, 61078002, 61107820, 11104214, 11081007, 11104216), and RFPD (20102210120005).

**Author contributions**

H.C. and Y.Z. wrote the main manuscript text and contributed to the theoretical and experimental analysis of this work. Y.X., Z.W. and X.Z. prepared figures 1–4. Y.Z. and M.X. provided the idea. All authors reviewed the manuscript.

**Additional information**

Competing financial interests: The authors declare no competing financial interests.

How to cite this article: Chen, H.X. et al. Parametrically Amplified Bright-state Polarization of Four- and Six-Wave Mixing in an Optical Ring Cavity. *Sci. Rep.* 4, 3619, DOI:10.1038/srep03619 (2014).

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