Higher-order squeezing for the codirectional Kerr nonlinear coupler

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In this Letter we study the evolution of the higher-order squeezing, namely, nth-order single-mode squeezing, sum- and difference-squeezing for the codirectional Kerr nonlinear coupler. We show that the amount of squeezing decreases when \( n \), i.e. the squeezing order, increases. For specific values of the interaction parameters squeezing factors exhibit a series of revival-collapse phenomena, which become more pronounced when the value of \( n \) increases. Sum-squeezing can provide amounts of squeezing greater than those produced by the nth higher-order \((n \geq 2)\) squeezing for the same values of interaction parameters and can map onto amplitude-squared squeezing. Further, we prove that the difference-squeezing is not relevant measure for obtaining information about squeezing from this device.

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I. INTRODUCTION

The optical coupler is a device composed of two (or more) waveguides, which are placed close enough to each others. The guided modes are coupled by means of the evanescent waves and hence the energy exchanged between the waveguides can be controlled [1]. In this regard directional coupler is an important device for data transmission and optical communication networks [2]. Such device has been experimentally implemented, e.g. in [3]. Directional coupler involving Kerr nonlinearity is an important device owing to its application in optics as an intensity-dependent routing switch [1, 4]. The quantum properties for this device have been studied by several authors [5, 6, 7, 8, 9, 10, 11, 12]. For more details the reader can
consult the review paper [13].

As is well known that squeezed light has less noise than coherent light in one of the field quadratures provided that the uncertainty relation is fulfilled. This light has various application, e.g. in quantum information, high precision measurements, etc. This encourages researchers for developing different types of squeezing. For instance, higher-order squeezing of a single-mode case was suggested and examined in [14]. In this direction the definitions for amplitude-squared squeezing [15], amplitude-cubed squeezing [16] and the \( n \)th power squeezing [17] have been developed. Furthermore, higher-order two-mode squeezing has been given in the sense of the sum- and difference-squeezing [18]. Actually, the term higher-order is given for the sum- and difference-squeezing since the quadrature operators are defined in terms of a product (not a sum) of mode operators. Quite recently, general multimode sum-squeezing [19] and difference-squeezing [20] have been adopted. Such definitions for higher-order squeezing are motivated by the development in the higher-order correlation measurement techniques aiming to extract information efficiently from the optical signal [14]. It is worth mentioning that the first experimentally observed squeezed states are of the two-mode type [21]. Also like one that exhibits second-order (normal) squeezing, a field that is squeezed to a higher order is a pure quantum mechanical light and has no classical description.

Generally, the earlier investigation given to CKNC has been entirely focused on the normal squeezing, e.g. [7, 13]. Thus in this Letter we investigate the evolution of the higher-order squeezing involving also the sum- and difference-squeezing. This will be done in the following order. In section 2 we give the basic relations and equations, which will be used in the paper. In section 3 we investigate and discuss the results. In section 4 we give the main conclusions.

\section{II. BASIC EQUATIONS AND RELATIONS}

In this section we give the basic equations and relations, which include the Hamiltonian formalism for the system under consideration, the solutions for the equations of motion and the general definition for squeezing.
The Hamiltonian for the codirectional Kerr nonlinear coupler (CKNC) is
\[
\frac{\hat{H}}{\hbar} = \sum_{j=1}^{2} [\omega_j \hat{a}_j^\dagger \hat{a}_j + \chi \hat{a}_j^\dagger^2 \hat{a}_j^2] + \bar{\chi} \hat{a}_1^\dagger \hat{a}_1 \hat{a}_2 + \kappa (\hat{a}_1^\dagger \hat{a}_2 + \hat{a}_2^\dagger \hat{a}_1),
\]
where \(\omega_1\) and \(\omega_2\) are the frequencies of the first and the second modes with the annihilation operators \(\hat{a}_1\) and \(\hat{a}_2\), respectively, \(\chi\) and \(\bar{\chi}\) are the coupling constants proportional to the third-order susceptibility \(\chi^{(3)}\) and responsible for the self-action and cross-action processes, respectively, \(\kappa\) is the linear coupling constant between the waveguides. The solution of the Heisenberg equations related to (1) when \(\bar{\chi} = 2\chi\), can be easily obtained as:
\[
\hat{a}_1(t) = \exp(-i\hat{\Lambda}t/2) \left\{ \hat{a}_1(0) \left[ \cos(\lambda t) - i \frac{\Delta}{2\chi} \sin(\lambda t) \right] - i \frac{\chi}{\bar{\chi}} \hat{a}_2(0) \sin(\lambda t) \right\},
\]
\[
\hat{a}_2(t) = \exp(-i\hat{\Lambda}t/2) \left\{ \hat{a}_2(0) \left[ \cos(\lambda t) - i \frac{\Delta}{2\chi} \sin(\lambda t) \right] - i \frac{\chi}{\bar{\chi}} \hat{a}_1(0) \sin(\lambda t) \right\},
\]
where \(\lambda = \sqrt{\kappa^2 + \frac{1}{4}\Delta^2}\), \(\hat{\Lambda} = \omega_1 + \omega_2 + 4\chi(\hat{a}_1^\dagger \hat{a}_1 + \hat{a}_2^\dagger \hat{a}_2)\) and \(\Delta\) is the frequency mismatch. It is obvious that \(\hat{a}_1(t) \leftrightarrow \hat{a}_2(t)\) when \(\hat{a}_1(0) \leftrightarrow \hat{a}_2(0)\). The nature of the coupler, i.e. the switching of energy between waveguides, manifests itself as periodic functions in (2). Moreover, in addition with the energy exchange, both optical fields in the CKNC undergo the self-phase modulation owing to nonlinearity in the waveguides described by the cubic susceptibility \(\chi^{(3)}\) and would manifest itself in the equations as a nonlinear-modulation phase term, as we shall see. Assuming that the two modes are initially prepared in the coherent light \(|\alpha_1, \alpha_2\rangle\), one can evaluate the general form for the different moments of the operators \(\hat{A}_j(t) = \hat{a}_j(t) \exp[\frac{i}{\hbar} \omega_j (t_1 + \omega_2)]\), where \(\hat{a}_j(t)\) are given by (2), as
\[
\langle \hat{A}^m_1(t) \hat{A}^n_2(t) \hat{A}^{m_1}_1(t) \hat{A}^{n_2}_2(t) \rangle = \exp[\{(|\alpha_1|^2 + |\alpha_2|^2) (z^{n_2+n_3-n_1-1})
\]
\[
\times \alpha_1^{m_2}(t) \alpha_2^{n_4}(t) \alpha_1^{s_1}(t) \alpha_2^{s_2}(t) z^{[n_2 n_4 + \frac{\alpha_1}{\chi} (n_2 - 1) + \frac{\alpha_2}{\chi} (n_4 - 1) - n_1 n_3 - \frac{\alpha_1}{\chi} (n_1 - 1) - \frac{\alpha_2}{\chi} (n_3 - 1)]}
\]
where \(n_j, j = 1, 2, 3, 4\) are integers,
\[
z = \exp(-2i\chi t), \quad \alpha_1(t) = \alpha_x(t) + i\alpha_y(t), \quad \alpha_2(t) = \alpha_x'(t) + i\alpha_y'(t)
\]
(4) and
\[
\alpha_x(t) = \alpha_1 \cos \lambda t, \quad \alpha_y(t) = -[\alpha_1 \frac{\Delta}{2} + \alpha_2 \kappa] \frac{\sin \lambda t}{\lambda},
\]
\[
\alpha_x'(t) = \alpha_2 \cos \lambda t, \quad \alpha_y'(t) = -[\alpha_2 \frac{\Delta}{2} + \alpha_1 \kappa] \frac{\sin \lambda t}{\lambda}.
\]
(5)
We have assumed that $\alpha_1$ and $\alpha_2$ are real.

On the other hand, for investigating squeezing we have to define two quadratures $\hat{X}$ and $\hat{Y}$, which denote the real (electric) and imaginary (magnetic) parts of the radiation field. Assume that these quadratures satisfy the following commutation rule:

$$[\hat{X}, \hat{Y}] = \frac{\hat{C}}{2},$$

where $\hat{C}$ may be $c$-number or operator. The uncertainty relation associated with the commutation rule (6) is

$$\langle (\Delta \hat{X})^2 \rangle \langle (\Delta \hat{Y})^2 \rangle \geq \frac{|\langle \hat{C} \rangle|^2}{16},$$

where $\langle (\Delta \hat{X})^2 \rangle = \langle \hat{X}^2 \rangle - \langle \hat{X} \rangle^2$ and similar form can be given for $\langle (\Delta \hat{Y})^2 \rangle$. The system is said to be squeezed in the $X$-quadrature if

$$S = 4\langle (\Delta \hat{X}(t))^2 \rangle - \frac{|\langle \hat{C} \rangle|}{|\langle \hat{C} \rangle|} \leq 0.$$  

The equality sign in (8) holds for minimum-uncertainty states. Similar definition can be given for the $Y$-quadrature (defining a $Q$-factor). Equations (3)–(8) provide all the necessary tools to describe evolution of the different types of higher-order squeezing.

III. RESULTS AND DISCUSSIONS

In this section we discuss three types of squeezing, which are $n$th-order single-mode squeezing, sum-squeezing and difference-squeezing. This will be performed in the following parts.

A. The $n$th-order single-mode squeezing

In this part we treat the $n$th-order single-mode squeezing. For convenience we use the definition given in [17]. In this case $\hat{X}, \hat{Y}$ and $\hat{C}$ take the forms:

$$\hat{X} = \frac{1}{2}[\hat{A}_1^n(t) + \hat{A}_1^{\dagger n}(t)], \quad \hat{Y} = \frac{1}{2i}[\hat{A}_1^n(t) - \hat{A}_1^{\dagger n}(t)], \quad \hat{C} = \hat{A}_1^n(t)\hat{A}_1^{\dagger n}(t) - \hat{A}_1^{\dagger n}(t)\hat{A}_1^n(t).$$
where \( n \) is positive integer. For obtaining some accurate information we assume that \( \alpha_1 = \alpha_2 = \alpha \) and \( \Delta = 0 \) (resonance case). From (3), (8) and (9) we can obtain

\[
S_1(t) = \mu [1 + h_1(t) - h_2(t)], \quad Q_1(t) = \mu [1 - h_1(t) - h_3(t)],
\]

where

\[
\begin{align*}
    h_1(t) &= \cos[2\lambda t n + 2n(2n - 1)\chi t + \epsilon \sin(4n\chi t)]f(2n\chi t), \\
    h_2(t) &= 2 \cos^2[\lambda t n + n(n - 1)\chi t + \epsilon \sin(2n\chi t)]f^2(n\chi t), \\
    h_3(t) &= 2 \sin^2[\lambda t n + n(n - 1)\chi t + \epsilon \sin(2n\chi t)]f^2(n\chi t), \\
    \epsilon &= |\alpha_1|^2 + |\alpha_2|^2, \quad \mu = \frac{2\alpha^2n}{\langle \hat{C} \rangle}, \\
    f(n\chi t) &= \exp[-2\epsilon \sin^2(n\chi t)].
\end{align*}
\]

Actually, the origin of occurrence of the nonclassical effects in CKNC is in the existence of the envelope function and/or of the nonlinear-modulation phase term \( f(n\chi t) \). The evolution of this function is mainly responsible for the features of the squeezing factors. In this regard the value of the parameter \( \chi t \) plays the crucial role in obtaining squeezing. The envelope function is periodic and the period decreases as the value of \( n \) increases. Roughly speaking, it is obvious that squeezing occurs when the amount in the brackets of (10)–which is finite–is less than zero. The pre-factor \( \mu \) plays an amplification role. Suppose that \( \chi t = \frac{\pi}{m} \). Thus for \( m = n \) or \( m = 1 \) the system reduces to its initial stage, i.e. disentangled coherent states but the amplitudes may be different from those of the initial ones. Consequently, squeezing may occur and switching between the two waveguides only when \( n/m = l, l \) is fraction. Now we prove that for specific values of interaction time the system can exhibit higher-order squeezing. For instance, for \( n \) odd, i.e. \( n(2n - 1) \) is odd and \( n(n - 1) \) is even, and \( \chi t = \pi/2 \), say, the expressions (10) reduces to

\[
\begin{align*}
    S_1(t) &= \mu \left\{1 - \cos(2n\lambda t) - 2 \cos^2(n\lambda t) \exp(-2\epsilon)\right\}, \\
    Q_1(t) &= \mu \left\{1 + \cos(2n\lambda t) - 2 \sin^2(n\lambda t) \exp(-2\epsilon)\right\}.
\end{align*}
\]
FIG. 1: Evolution of the squeezing factor $S_1(t)$ of the first mode when $\kappa = 1, \chi = 0.5, (\alpha_1, \alpha_2) = (2, 0)$ and for (a) $\Delta = 0$ (solid curve for $n = 2$ and dashed curve for $n = 3$), (b) $(\Delta, n) = (50, 2)$, and (c) $(\Delta, n) = (50, 3)$.

Therefore, for $\lambda t = \pi/n$, say, squeezing can be only observed in $S_1(t)$, whereas for $\lambda t = \pi/(2n)$ it is only obtained in $Q_1(t)$. One can realize that the values of interaction time to which squeezing occur, depend on the order of the squeezing. On the other hand, when $n$ is even and $\chi t = \pi/2$ one can prove that the system provides its initial stage, i.e. disentangled coherent states. As we mentioned above the quantity in the brackets of (10) is finite, i.e. its value locates in the intervals $[0, \pm 2]$, so the natural question is that which value of $n$ provides maximum squeezing? The answer for this can be obtained by examining the amplification factor $\mu$ for different values of $n$. We found that as $n$ increases the value of $\mu$ decreases. This means that the best value of squeezing can be obtained for the lowest order, i.e. $n = 1$ for normal squeezing [7]. Nevertheless, the periodicity of occurring squeezing in the time domain increases as $n$ increases.

These facts are remarkable in Figs. (1) and (2) for strong and weak intensities, respectively, for given values of interaction parameters. From Fig. 1(a) one can observe that squeezing occurs periodically and the periodicity increases as well as the amount of squeezing decreases when the order $n$ increases (compare the dashed and solid curves). The value of the detuning parameter $\Delta$ plays an important role (see Figs. 1(b) and (c)). From Figs. 1(b) and (c) for $\Delta >> 1$, $S_1(t)$ exhibits particular shape of periodic revival-collapse phenomenon and the amount of the nonclassical squeezing becomes much more pronounced than before (compare Fig. 1(b) and (c) to the solid and dashed curves in Fig. 1(a), re-
FIG. 2: Evolution of the squeezing factor $S_1(t)$ of the first mode when $\kappa = 1, \chi = 0.5, (\alpha_1, \alpha_2) = (0.3, 0.3), \Delta = 50$, and for (a) $n = 2$, (b) $n = 3$.

spectively. Generally, the occurrence of revival-collapse phenomenon in the evolution of squeezing factors may be explained as follows. Basically $S_1(t)$ includes two forms of periodic function, namely, the trigonometric and envelope functions. These functions are periodic in the two parameters $\lambda t$ (with period $\pi/n\lambda$) and $\chi t$ (with period $\pi/n\chi$). Generally, when the values of $\Delta$ increase, the period of the energy exchange between waveguides decreases, i.e. many oscillations occur, till the interaction time becomes $t\chi = \pi/2n$, at this moment the field is trapped instantaneously by nonlinearity in the waveguides and the squeezing factors show collapse. As the interaction proceeds the phenomenon is periodically repeated. The sensitivity of the revival-collapse phenomenon to the value of $n$ can be realized by comparing Fig. 1(b) and Fig. 1(c). The number of the revival patterns increases as $n$ increases. The particular shape of revival-collapse phenomenon in Figs. 1(b) and (c) can be understood as follows. In (10) we have two forms of the envelope function, which are $f(2n\chi t)$ and $f(n\chi t)$. As we mentioned above these are periodic functions, however, the period of the first one is two times less than that of the second. Consequently, for $\chi t = \frac{m\pi}{n}, m$ is integer, $h_1(t)$ and $h_2(t)$ provide simultaneously their maximum contribution, which interfere with each others producing squeezing in $S_1(t)$, whereas for $\chi t = \frac{m\pi}{n}, m$ is odd integer, $h_1(t)$ provides the main contribution and hence complete revivals occur. Such behavior can be realized only for large
intensities, where $f(n\chi t) = 0$ or $1$, however, for weak intensities the behavior will be rather different. This is related to the fact that $f(n\chi t) \simeq 1$ everywhere (compare Figs. (1) and (2)). Now we draw the attention to Figs. 2(a) and (b). Comparison between Figs. 1 and 2 leads to the result that the value of squeezing for strong-intensity regime is much greater than that for the weak-intensity regime and the shapes of revival-collapse phenomenon in the two regimes are quite different. The origin of this difference is that for weak intensities there is no exact collapse causing that the revival patterns are much broader than that of the strong-intensity regime. Also the various facts mentioned above related to the value of $n$ are still valid in the weak-intensity regime. Throughout the discussion we have focused on the evolution of $S_1(t)$ because we have noted that for $\Delta = 0$, $Q_1(t)$ is almost positive, i.e. it cannot provide squeezing, and for $\Delta \neq 0$ it provides typical forms as those for the $S_1(t)$ (see Figs. 1(b), (c) and Figs. 2).

B. Sum-squeezing

It is worth reminding that the sum- and difference-squeezing has been realized in nonlinear optics for four-wave sum [22] and difference [23] frequency generation. We proceed by using the definition given in [18] for sum-squeezing. In this case the operators $\hat{X}$, $\hat{Y}$ and $\hat{C}$ take the forms:

$$\hat{X} = \frac{1}{2}[\hat{A}_1(t)\hat{A}_2(t)+\hat{A}_{1}^\dagger(t)\hat{A}_{2}^\dagger(t)], \quad \hat{Y} = \frac{1}{2t}[\hat{A}_1(t)\hat{A}_2(t)-\hat{A}_{1}^\dagger(t)\hat{A}_{2}^\dagger(t)], \quad \hat{C} = \hat{N}_1+\hat{N}_2+1. \quad (13)$$

One can easily check that when $\alpha_1 = \alpha_2$ sum-squeezing map onto the amplitude-squared squeezing given above, i.e. $n = 2$ (see Figs. 1). Here we pay attention to the case $\alpha_1 \neq \alpha_2$. Assume that $\alpha_1 = \alpha, \alpha_2 = 0$ and $\Delta = 0$. For this case the sum-squeezing factors can be
evaluated as

\[
S_2(t) = \frac{2a^4}{\alpha^2+1} \sin^2(\lambda t) \cos^2(\lambda t) \left\{ 1 + \cos(12\chi t + \epsilon \sin(8\chi t)) \exp[-2\epsilon \sin^2(4\chi t)] \right. \\
- 2\sin^2(2\chi t + \epsilon \sin(4\chi t)) \exp[-4\epsilon \sin^2(2\chi t)] \left. \right\},
\]

\[
Q_2(t) = \frac{2a^4}{\alpha^2+1} \sin^2(\lambda t) \cos^2(\lambda t) \left\{ 1 - \cos(12\chi t + \epsilon \sin(8\chi t)) \exp[-2\epsilon \sin^2(4\chi t)] \right. \\
- 2\cos^2(2\chi t + \epsilon \sin(4\chi t)) \exp[-4\epsilon \sin^2(2\chi t)] \left. \right\}.
\]

Expressions (14) show that the system is able to produce sum-squeezing, e.g. when \(\chi t = \pi/4\) these expressions reduce to

\[
S_2(t) = -\frac{a^4}{\alpha^2+1} \sin^2(2\lambda t) \exp(-2\epsilon), \\
Q_2(t) = \frac{a^4}{\alpha^2+1} \sin^2(2\lambda t).
\]

It is evident that squeezing can occur in the first quadrature. Figs. 3(a) and (b) are given for \(S_2(t)\) for strong and weak intensity regimes, respectively, for the given values of the parameters. One can observe that squeezing occurs periodically and becomes more pronounced when the intensities increase (see the inset in Fig. 3(b), also compare Fig. 3(a) and Fig. 3(b)). Influence of the detuning parameter is shown in Fig. 3(b), which manifests itself as the revival-collapse phenomenon. Explanation as that given in the first part can be given here. Further, we have noted that the nonclassical values of the sum-squeezing are much greater than those for the \(n\)th-order single-mode squeezing \((n \geq 2)\) for the same values of the interaction parameters. Also for \(\Delta \gg 1, \alpha_j > 1\) Figs. 1(b) and (c) are obtained. Finally, conclusions similar to those given for \(Q_1(t)\) in the first part are valid for \(Q_2(t)\).

C. Difference-squeezing

In this part we show that difference-squeezing factors fail to give information about squeezing from the coupler. For difference-squeezing the operators \(\hat{X}, \hat{Y}\) and \(\hat{C}\) take the forms [18]:

\[
\hat{X} = \frac{1}{2}[\hat{A}_1(t)\hat{A}_2^\dagger(t) + \hat{A}_1^\dagger(t)\hat{A}(t)], \\
\hat{Y} = \frac{1}{2i}[\hat{A}_1(t)\hat{A}_2^\dagger(t) - \hat{A}_1^\dagger(t)\hat{A}(t)], \\
\hat{C} = \hat{N}_2 - \hat{N}_1.
\]
From (3), (8) and (16) we can obtain the difference-squeezing factors as

\[ S_3(t) = 2\text{Re}[\bar{\alpha}_1^2(t)\bar{\alpha}_2^2(t)] + 2|\bar{\alpha}_1(t)|^2|\bar{\alpha}_2(t)|^2 + 2|\bar{\alpha}_1(t)|^2 - [\bar{\alpha}_x(t)\bar{\alpha}_x'(t) + \bar{\alpha}_y(t)\bar{\alpha}_y'(t)]^2, \]

(17)

\[ Q_3(t) = -2\text{Re}[\bar{\alpha}_1^2(t)\bar{\alpha}_2^2(t)] + 2|\bar{\alpha}_1(t)|^2|\bar{\alpha}_2(t)|^2 + 2|\bar{\alpha}_1(t)|^2 - [\bar{\alpha}_x(t)\bar{\alpha}_x'(t) - \bar{\alpha}_x(t)\bar{\alpha}_x'(t)]^2, \]

where \( \bar{\alpha}_j(t) \) are given by (11) and (15) and \( \text{Re} \) stands for real value. It is evident that (17) is independent of the nonlinear-modulation phase term and consequently the system cannot provide difference-squeezing. This can be confirmed after minor manipulation with (17), which reduces to

\[ S_3(t) = Q_3(t) = 2|\bar{\alpha}_1(t)|^2. \]

(18)

This means that the two squeezing factors are typical and equals twice the mean photon number in the first waveguide. Such behavior of difference-squeezing can be understood by noting that the coupler and the quadratures of the difference-squeezing are describing by the same mechanism. To be more specific, the quadratures of the difference-squeezing represent up conversion processes (cf. (16)), i.e. when one photon is created in the first mode the other is annihilated in the second mode. On the other hand, the coupler is basically operating by switching energy between waveguides (conservation of energy).

IV. CONCLUSIONS

Throughout this Letter we have studied for the first time the higher-order squeezing for CKNC. For the \( n \)th-order single-mode squeezing we have found that the amount of squeezing decreases as the order of the squeezing increases regardless of the values of the intensities. Frequency mismatch can increase (or generate) squeezing in the quadratures. Also squeezing factors exhibit revival-collapse phenomenon resulting from the competition between the Kerr nonlinearity and the frequency mismatch. Further, the number of the revival patterns increase when the order of squeezing increases. The locations of these revival patterns in the time domain depend on the values of \( nt\chi \), whereas their shapes depend on the intensities of the field launched in the waveguides initially. Furthermore, sum-squeezing can map onto amplitude-squared squeezing when the intensities are equal and can provide revival-collapse phenomenon based on the values of \( \Delta \). Sum-squeezing can provide amounts of squeezing greater than those produced by the \( n \)th higher-order \((n \geq 2)\) squeezing for
FIG. 3: Evolution of the sum-squeezing factor $S_2(t)$ for $\kappa = 1, \chi = 0.5$, and for (a) $(\Delta, \alpha_1, \alpha_2) = (0, 1, 1.5)$ (short-dashed curve) and $(0, 2, 3)$ (solid curve), (b) $(\Delta, \alpha_1, \alpha_2) = (50, 0.3, 0.6)$. The inset in (a) is given for the sake of comparison.

the same values of interaction parameters. This means that the sum-squeezing is a better measure for extracting information about squeezing from CKNC. These conclusions are in relation to the structure of the nonlinear part of the Hamiltonian (1). Also we have proved that the difference-squeezing is not suitable for extracting information about squeezing from CKNC. The final remark, we have numerically noted that the occurrence of the revival-collapse phenomenon in the squeezing factors depends on the value of $\Delta$ and not on $\kappa$, i.e. on the intensity of the linear exchange between waveguides.

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