

Nondegenerate bound-state solitons in multicomponent Bose-Einstein condensates

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to observe them are discussed in BEC systems. Finally, we summarize our results in Sec. VI.

II. THEORETICAL MODEL AND NONDEGENERATE BOUND-STATE SOLITON SOLUTIONS

In the framework of the mean-field theory, the dynamics of two-component BECs can be described well by the following dimensionless two-component coupled nonlinear Schrödinger equations [25]:

$$\begin{aligned} iq_{1,t} + q_{1,xx} + 2(|q_1|^2 + |q_2|^2)q_1 &= 0, \\ iq_{2,t} + q_{2,xx} + 2(|q_1|^2 + |q_2|^2)q_2 &= 0, \end{aligned} \quad (1)$$

where $q_1(x, t)$ and $q_2(x, t)$ represent the mean-field wave functions of the two-component BECs [26]. The interactions between atoms are attractive for the above model, and similar discussions can be had for repulsive interactions cases. The model can be also used to describe the evolution of light in nonlinear optical fibers [27,28]. With the aid of the Darboux transformation (DT) [29–32], the Hirota method [33,34], the Bäcklund transformation [35], and the Kakomtsev-Petviashvili hierarchy reduction method [36–38], many different vector solitons have been obtained in various nonlinear systems, such as bright-bright solitons [2,39,40], bright-dark solitons [5], dark-dark solitons [7,8], and dark-bright solitons [9,10,35]. We would like to look for NDBSSs, for which all eigenstates are bound states, in contrast to the reported bright-dark soliton [5] and the dark-bright soliton [9,10,35].

DT is a very powerful method for constructing soliton solutions. Therefore, we would like to develop a DT method to derive the NDBSSs. The twofold DT with spectral parameters $\lambda_1 = a_1 + ib_1$ and $\lambda_2 = a_1 + ib_2$ generates one NDBSS. The derivation process is given in Appendix A in detail, which is different from the processes for generating bright-bright solitons and bright-dark solitons [29–32]. The exact general double-hump soliton solution for Eq. (1) can be written as follows:

$$\begin{aligned} q_1(x, t) &= -2ib_1c_{11}^* \frac{N_1}{M_1} e^{-i[a_1x + (a_1^2 - b_1^2)t]}, \\ q_2(x, t) &= -2ib_2c_{22}^* \frac{N_2}{M_1} e^{-i[a_1x + (a_1^2 - b_2^2)t]}. \end{aligned} \quad (2)$$

with

$$\begin{aligned} N_1 &= \left(\frac{b_1 - b_2}{b_1 + b_2} + |c_{22}|^2 e^{-2b_2(x+2a_1t)} \right) e^{-b_1(x+2a_1t)}, \\ N_2 &= \left(\frac{b_2 - b_1}{b_1 + b_2} + |c_{11}|^2 e^{-2b_1(x+2a_1t)} \right) e^{-b_2(x+2a_1t)}, \\ M_1 &= |c_{11}|^2 e^{-2b_1(x+2a_1t)} + |c_{22}|^2 e^{-2b_2(x+2a_1t)} \\ &\quad + |c_{11}c_{22}|^2 e^{-2(b_1+b_2)(x+2a_1t)} + \frac{(b_1 - b_2)^2}{(b_1 + b_2)^2}, \end{aligned}$$

where a_1 , b_1 , and b_2 are real parameters, and c_{11} and c_{22} are complex parameters. The velocity of the soliton is $-2a_1$. The parameters b_1 , b_2 , $|c_{11}|$, and $|c_{22}|$ determine the soliton profile. According to the expressions in Eq. (2), we see that soliton profiles will be kept except for the translation of the position under the transforms $|c_{11}| \rightarrow |c_{11}|e^{b_1\delta}$ and $|c_{22}| \rightarrow |c_{22}|e^{b_2\delta}$,

where δ is an arbitrary real constant. These parameters non-trivially contribute to the profile of the soliton. It is noted that the solitons in two components admit different modes. When $b_1 < b_2$, the soliton in the q_2 component admits no node, and the one in the q_1 component always has one node, and vice versa. Since a nonlinear Schrödinger equation can be mapped to a linear Schrödinger equation with proper quantum wells [11,12], the nonlinear term $-2|q_1(x, t)|^2 - 2|q_2(x, t)|^2$ in Eq. (1) can be understood as an effective quantum well, $V(x, t)$. If $-2|q_1(x, t)|^2 - 2|q_2(x, t)|^2$ does not depend on time, the soliton states can be related with eigenproblems in the quantum well $V(x)$, namely, $[-\partial_x^2 + V(x)]q_j = \mu_j q_j$. Then the soliton solutions can be mapped to the eigenvalues μ_j and the eigenstates q_j in the effective quantum well. From the node properties of the corresponding eigenstates in quantum wells [12], we know that the eigenstate with one node corresponds to the first-excited state in effective quantum wells [11]. Therefore, the bound-state solitons in two components correspond to ground state and first-excited state in effective quantum wells respectively. A bound-state soliton with one node is in sharp contrast to the bright soliton and the dark soliton reported before. The NDBSSs are similar to the “soliton complexes” which have been used to describe different composition solitons [41–43]; however, it is not convenient to distinguish nodes of solitons based on the concept of “soliton complexes.” To emphasize the node properties and their relations with eigenstates in quantum wells, we discuss them from the quantum well viewpoint. Moreover, the Manakov model can be transformed into the parametrically driven nonlinear Schrödinger equation [42,43], when the solutions can be transformed into stationary states of the Manakov model. In this way, NDBSSs can be related with the soliton complexes in Refs. [42,43]. Additionally, for the soliton solutions which cannot be transformed into stationary states of the Manakov model, they do not hold anymore for the parametrically driven nonlinear Schrödinger equation. In what follows, we discuss the profile types of the fundamental NDBSSs.

The profiles of NDBSSs can be mainly classified as three different types: asymmetric double-hump solitons (two components both admit asymmetric double-hump solitons), symmetric single-hump–double-hump solitons (one component admits single-hump solitons and the other component has a symmetric double-hump soliton), and symmetric double-hump solitons (two components both admit symmetric double-hump solitons). This classification is different from the one given in Ref. [44], based on different aspects of soliton profiles. The three different cases are shown in Figs. 1(a1)–1(a3); the solid blue line and the dashed red line corresponds to the q_1 component and the q_2 component, respectively. Figure. 1(a1) depicts the intensity profiles of asymmetric double-hump solitons in both components. The solitons in the q_1 component and the q_2 component correspond to the first-excited state and the ground state, respectively, in an effective quantum double-well potential. Particularly, we find that the asymmetric double-hump soliton solution, Eq. (2), can be reduced to a symmetric form by choosing the parameters $c_{11} = c_{22} = \sqrt{3}/3$, $b_2 = 2b_1$, and $\delta = 0$. For this case, we rewrite the solutions as $q_1 = \sqrt{3}b_1 \operatorname{sech}[b_1(x + 2a_1t)] \tanh[b_1(x + 2a_1t)] e^{-i[a_1x + (a_1^2 - b_1^2)t - \pi/2]}$ and $q_2 = \sqrt{3}b_1$

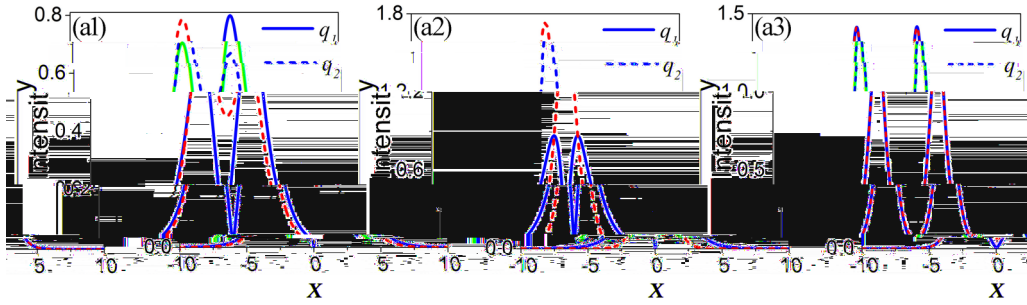


FIG. 1. Three different density profiles for double-hump solitons in two-component coupled systems: (a1) asymmetric double-hump solitons in both components, (a2) symmetric single-hump-double-hump solitons, and (a3) approximate symmetric double-hump solitons in both components. solid blue line and the dashed red line correspond to the q_1 component and the q_2 component, respectively. It is seen that solitons in q_1 and q_2 components correspond to first-excited state and ground state, respectively, in the effective double-well potential. The parameters are as follows: (a1) $c_{11} = c_{22} = 1$, $a_1 = 0$, $b_1 = 1$, $b_2 = 1.1$, and $\delta = -2.8$; (a2) $c_{11} = c_{22} = \sqrt{3}/3$, $a_1 = 0$, $b_1 = 1$, $b_2 = 2b_1$, and $\delta = 0$; and (a3) $c_{11} = c_{22} = 1$, $a_1 = 0$, $b_1 = 2$, $b_2 = 2.001$, and $\delta = -4.1$.

$\text{sech}^2[b_1(x + 2a_1t)]e^{-i[a_1x + (a_1^2 - 4b_1^2)t + \pi/2]}$. The amplitudes of the q_1 and q_2 components are $\sqrt{3}b_1/2$ and $\sqrt{3}b_1$, respectively. One can see that a symmetric double-hump bright soliton presents in the q_1 component, which corresponds to the first-excited bound state, while a single-hump ground bound-state soliton emerges in the q_2 component. This type of soliton can be seen as a symmetric single-hump-double-hump soliton. As an example, we show it in Fig. 1(a2). Additionally, when b_1 and b_2 are very close to each other, Eq. (2) will show nearly symmetrical double-hump bright solitons in both components. Typical intensity profiles are shown in Fig. 1(a3). The two humps distribute symmetrically in each component for this case. Remarkably, it is clear that a bright soliton with one node (the first-excited bound-state soliton) emerges in just one of the components for all three cases.

The soliton solutions of the above Manakov model have been studied widely. However, bright solitons with nodes are absent in most of the previously published studies on vector solitons [29–32]. Therefore, we would like to discuss why there are NDBSSs in the coupled systems. Based on the derivation method for NDBSSs, we note that these bound-state solitons are generated from two solitons with zero relative velocity and some special parameters. The analysis indicates that the choice of special parameters corresponds to the incoherent interactions between solitons in different components. The incoherent interaction corresponds to the nonlinear interaction between incoherent solitons, for which the relative phase of solitons does not affect the interaction pattern. It is different from the coherent interaction of two solitons in the same component for which the relative phase of them affects the interaction pattern. The double-hump soliton is related to two incoherent solitons, for which each bright soliton emerges in just one certain component (they are phase separated), in contrast to the well-known bright-bright soliton. Therefore, we investigate the interaction of two incoherent solitons with different velocities, to investigate the role of relative velocity in forming the NDBSSs. This can be done exactly by changing the spectral parameters $\lambda_1 = a_1 + ib_1$ and $\lambda_2 = a_2 + ib_2$ in the solution Eq. (A5). We change the velocity of the soliton ($v_j = -2a_j$) in each component, and other parameters are fixed as in Fig. 1(a1). The relevant dy-

namical processes of incoherent interactions between solitons are depicted in Fig. 2, and the relative velocities of solitons ($v_2 - v_1$) in two components correspond to 0.08, 0.004, and 0.0004 from top to bottom, respectively (see captions for detailed parameter settings). The characters of incoherent interactions between solitons are shown clearly in Fig. 2(a). The incoherent interaction mainly includes two cases: incoherent collision and incoherent superposition. Incoherent superposition can be approached from incoherent collision by decreasing the relative velocity. One can see that the relative velocity between solitons is smaller, and the nonlinear

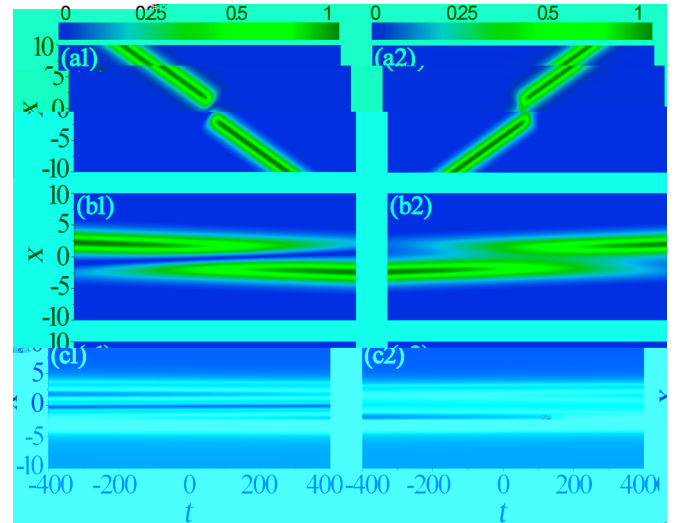


FIG. 2. The incoherent interactions between solitons in different components with different relative velocities ($2a_1 - 2a_2$). From top to bottom, the parameters are as follows: (a1) and (a2) $a_1 = -a_2 = 0.02$, (b1) and (b2) $a_1 = -a_2 = 0.001$, and (c1) and (c2) $a_1 = -a_2 = 0.0001$. The left panels show the density distributions of the q_1 component; the right panels show the density distributions of the q_2 component. The evolution dynamics suggest that nondegenerate bound-state solitons are induced by incoherent superpositions of bright solitons in different components. The other parameters are $c_{11} = c_{22} = 1$, $b_1 = 1$, $b_2 = 1.1$, and $\delta = -2.8$ [the same as those in Fig. 1(a1)].

incoherent interaction between solitons in different components is becoming stronger. When the relative velocity of solitons decreases to zero, the nonlinear incoherent interaction will become an incoherent superposition of them, which induces the profile of the bright soliton in one of component changes to be a double-hump soliton with one node (first-excited bound eigenstate) and the profile of the bright soliton in the other component becomes a double-hump soliton with no node (ground bound state), such as solitons in Fig. 1(a1). These dynamical processes indicate that NDBSSs are induced by the incoherent superposition between solitons in different components.

It should be mentioned that similar multihump solitons have been observed for a long time in dispersive nonlinear medium [45–47]. Stationary multihump solitons have also been obtained in Kerr-like media [48]. Recently, many different static nondegenerate solitons were derived from eigenstates in some certain quantum wells [11]. However, bright solitons with nodes are symbiotic with dark solitons; that is, the bound state and the free state always coexist in the coupled systems. Moreover, these solutions are stationary and it is inconvenient to discuss their collision properties analytically. Very recently, nondegenerate solitons in the Manakov system were reported by the Hirota's bilinear method [44]. In this paper, we develop the DT method to derive NDBSSs, which enables us to discuss the underlying mechanism for these bound-state solitons. Furthermore, we develop the DT method to investigate the collision processes between the NDBSSs analytically.

III. COLLISION BETWEEN DIFFERENT NONDEGENERATED SOLITONS

First, we investigate the collision between one NDBSS and one degenerate bright soliton, by performing the threefold DT with spectral parameters $\lambda_1 = a_1 + ib_1$, $\lambda_2 = a_1 + ib_2$, and $\lambda_3 = a_2 + ib_3$ (see Appendix A for the solution process). More complicated interactions between solitons can be investigated by performing the N -fold DT in Appendix A, Eq. (A8). For this case, typical densities are depicted in the Figs. 3(a1) and 3(a2) corresponding to the q_1 component and the q_2 component, respectively. It is seen that the interference patterns between an asymmetric double-hump soliton and a single-hump soliton appear in both components. One first-excited-state soliton interferes with a ground-state soliton in the q_1 component, while two ground-state solitons collide with each other in the q_2 component. Detailed analysis indicates that the collisions between them are usually inelastic, and they can be elastic under some special initial conditions.

Second, we investigate the interaction between two NDBSSs by performing the fourfold DT with spectral parameters $\lambda_1 = a_1 + ib_1$, $\lambda_2 = a_1 + ib_2$ (generates one NDBSS), $\lambda_3 = a_2 + ib_3$, and $\lambda_4 = a_2 + ib_4$ (generates another NDBSS). We exhibit the dynamical evolution of them in the right-hand panels of Fig. 3, based on the two double-hump solitons solution Eq. (A7). Figures 3(b1) and 3(b2) correspond to the q_1 component and the q_2 component, respectively. As one can see in Figs. 3(b1) and 3(b2), the collision of two identical symmetric double-hump solitons (two first-excited bound-state solitons) in the q_1 component and two identical single-hump solitons

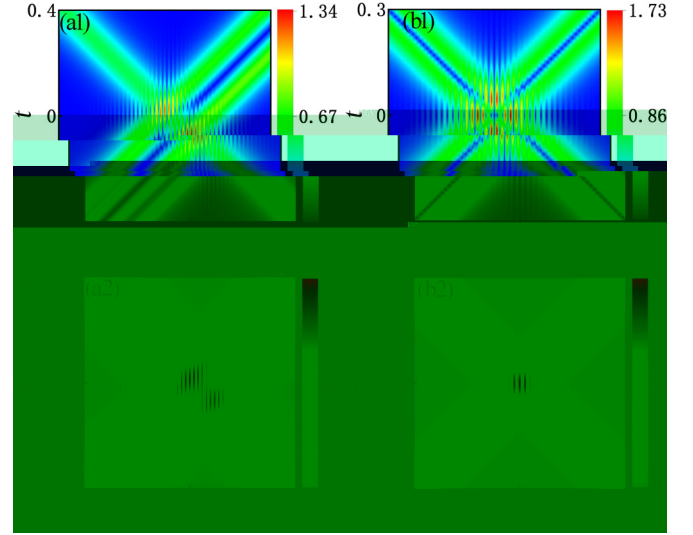


FIG. 3. The collision dynamics between one nondegenerate bound-state soliton and one degenerate bright soliton or one nondegenerate bound-state soliton. (a1) and (a2): The interference behavior between an asymmetric double-hump soliton (moving to the right direction) and a degenerate bright soliton (moving to the left direction). Related parameters are $c_{21} = c_{12} = 0$, $c_{11} = c_{22} = c_{13} = c_{23} = 1$, $a_1 = -10$, $a_2 = 10$, $b_1 = 1$, $b_2 = 1.1$, $b_3 = 0.8$, and $\delta_1 = \delta_2 = 0$. (b1) and (b2): The interference patterns between two nondegenerate bound-state solitons. The parameters are $c_{11} = c_{22} = c_{13} = c_{24} = \sqrt{3}/3$, $a_1 = 10$, $a_2 = -10$, $b_1 = 1$, $b_2 = 2$, $b_3 = 1$, $b_4 = 2$, and $\delta_1 = \delta_2 = 0$. The top panels show the density of the first component, and the bottom panels show the density of the second component. The soliton profiles change too slightly to be visible after the collision at these parameters.

(two ground bound-state solitons) in the q_2 component both generate interference patterns. Moreover, we further explore the interference properties of them by the asymptotic analysis technique (see the detailed solution process in Appendix A). Interestingly, we find that the interference of double-hump solitons presents multiperiodicity. The periodic functions are governed by the factors $\sin[(a_1 - a_2)x + (a_1^2 - a_2^2 + b_3^2 - b_1^2)t]$ and $\sin[(a_1 - a_2)x + (a_1^2 - a_2^2 - b_2^2 + b_4^2)t]$ and their corresponding cosine forms. These mean that there are three periodic oscillation behaviors in the interference process of two double-hump bright solitons. The spatial period is $D = 2\pi/(a_1 - a_2) = 4\pi/(v_1 - v_2)$, and the temporal periods are $T_1 = 2\pi/(a_1^2 - a_2^2 + b_3^2 - b_1^2)$ and $T_2 = 2\pi/(a_1^2 - a_2^2 - b_2^2 + b_4^2)$, in a striking contrast to interference patterns between bright solitons reported before [49]. This comes from more than two energy eigenvalues involved in the interference processes. In Figs. 3(b1) and 3(b2), only the spatial interference pattern is visible, due to the parameters settings $a_1^2 = a_2^2$, $b_1 = b_3$, and $b_2 = b_4$, which make two temporal periods be zero. For Figs. 3(a1) and 3(a2), the parameter settings make two temporal periods too small to be visible (see caption for details).

It should be emphasized that the collision dynamics between bound-state solitons in Fig. 3 are indeed inelastic, where the solitons' profiles change too slightly to be visible after the collision. Further analysis suggests that the collisions

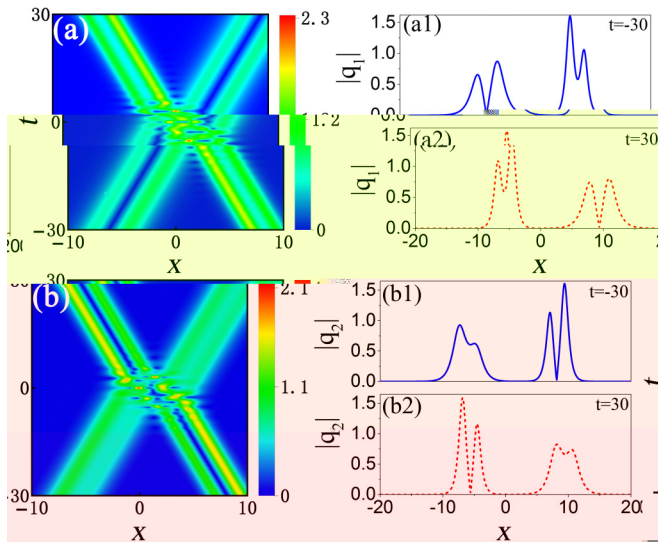


FIG. 4. The inelastic collision between two nondegenerate bound-state solitons. Left panels: Dynamical evolution for coherent collision of nondegenerate bound-state solitons. Right panels: Intensity plots for both components before ($t = -30$, solid blue line) and after ($t = 30$, dashed red line) the collision. It is seen that profiles of double-hump solitons indeed change after the collision. The analysis suggests that the collisions between these bound-state solitons are usually inelastic, due to the incoherent interaction and the coherent collision between them. The parameters are $c_{11} = c_{22} = c_{13} = c_{24} = 1$, $a_1 = -1/10$, $a_2 = 1/10$, $b_1 = 1$, $b_2 = 1.2$, $b_3 = 2$, $b_4 = 1.9$, and $\delta_1 = \delta_2 = 0$.

between bound-state solitons are usually inelastic unless the parameters satisfy the sufficient condition Eq. (A15). Inelastic collision refers to that the soliton's profile changes after the collision. A typical example for inelastic collision is shown in Fig. 4, where two double-hump solitons collide with each other in both components. The panels on the left show the related density evolutions; Figs. 4(a) and 4(b) correspond to the q_1 component and the q_2 component, respectively. The panels on the right depict the corresponding intensity profiles of two solitons before ($t = -30$, solid blue line) and after ($t = 30$, dashed red line) the collision in both components.

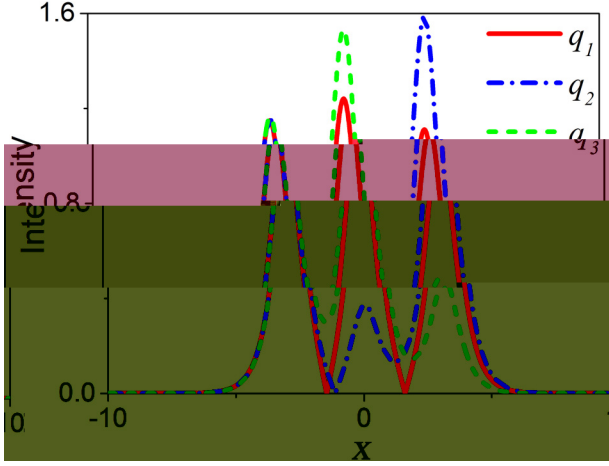


FIG. 5. Intensity profiles for triple-hump bright solitons of three-component BECs. The triple-hump bright soliton with no node in the q_3 component is the ground bound state (dashed green line), the triple-hump bright soliton with one node in the q_2 component is the first-excited bound state (dotted-dashed blue line), and the triple-hump bright soliton with two nodes in the q_1 component is the second-excited bound state (solid red line) in the effective quantum wells. The parameters are $c_1 = c_2 = c_3 = 1$, $a_1 = 0$, $b_1 = 1.99$, $b_2 = 2$, $b_3 = 2.02$, and $\delta = -5.2$.

solitons can be studied by performing the fivefold DT. The interplay between two triple-hump solitons can be explored by performing the sixfold DT. The inelastic collision of these bound-state solitons can be also expected.

V. STABILITY OF NONDEGENERATE BOUND-STATE SOLITONS

We discuss the stability of these bound-state solitons by calculating the Bogoliubov-de Gennes excitation spectrum around a stationary soliton solution. We introduce the ansatz $\mathbf{q} = \mathbf{q}_0 + \epsilon[\mathbf{P}e^{-i\lambda t} + \mathbf{Q}^*e^{i\lambda^* t}]e^{-i\mu t}$, where $\mathbf{q}_0$ is a NDBSS solution, ϵ is a small parameter, and $\{\lambda, (\mathbf{P}, \mathbf{Q}^*)\}$ defines an eigenvalue-eigenvector pair. Then, substituting this ansatz into related dynamical equations and linearizing the equations, we arrive at $O(\epsilon)$ at an eigenvalue problem for eigenvectors $(\mathbf{P}, \mathbf{Q}^*)$ and eigenvalues λ . Note that the latter may, in principle, be complex, i.e., $\lambda = \lambda_r + i\lambda_i$, instability corresponds to $\lambda_i \neq 0$. As an example, we consider the case of the asymmetric double-hump solitons in both components based on the solution Eq. (2). The corresponding eigenvalue spectrum is shown in Fig. 6, and the parameters are the same as those in Fig. 1(a). It can be seen that the spectrum of the double-hump solitons consists of purely real eigenvalues, showing spectral stability. Numerical simulations from the corresponding initial conditions also support the stability of solitons. On the other hand, experimental observations demonstrated that two-component solitons can be produced based on well-developed density and phase modulation techniques [9,51,52]. Solitons' interactions have also been demonstrated experimentally in BEC systems [53–56]. Therefore, based on the established experimental techniques,

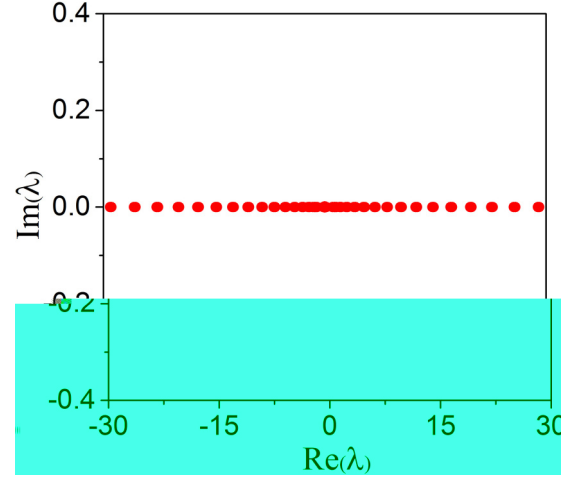


FIG. 6. The excitation spectrum of nondegenerate bound-state solitons. It is seen that nondegenerate bound-state solitons admit spectral stability. The related parameters are same as those in Fig. 1(a).

the presented NDBSSs have many possibilities to be observed in multicomponent attractive BECs.

VI. CONCLUSION AND DISCUSSION

In summary, we derive and investigate double-hump and triple-hump bound-state solitons in multicomponent BECs. The analysis indicates that bright solitons with nodes correspond to the excited bound eigenstates in the effective quantum wells. Particularly, we reveal that the incoherent superposition between solitons in different components is the generation mechanism of the bound-state solitons. Furthermore, we demonstrate that the collisions of NDBSSs are inelastic in general cases, which are induced by incoherent interactions and coherent collisions. More abundant bound-state solitons are expected in multi-component coupled BECs.

In fact, these NDBSSs can be also obtained from more general multisolitons (obtained using different methods) with some special constraint conditions. But the constraint conditions were not addressed in previous studies [34–36]. Our incoherent superposition form can be used to determine the constraint conditions. Moreover, NDBSSs can be mapped to eigenstates in a quantum well, in sharp contrast to the general multisolitons.

Note added. Recently, we noticed nondegenerate solitons were discussed in nonlinear optical fibers by the Hirota bilinear method [44]. In this paper, we perform the Darboux transformation method to derive NDBSS solutions. Moreover, the discussions on the mechanism and the node properties could be helpful for our understanding of NDBSSs.

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APPENDIX A: DERIVATION OF THE VECTOR DOUBLE-HUMP BRIGHT SOLITON SOLUTIONS OF EQ. (1)

The two-component coupled nonlinear Schrödinger equation, Eq. (1), is the compatibility condition of the linear spectral problems [25,31]:

$$\begin{aligned}\Phi_x &= U(x, t; \lambda)\Phi, \\ \Phi_t &= V(x, t; \lambda)\Phi,\end{aligned}\quad (\text{A1})$$

where

$$\begin{aligned}U &= \begin{pmatrix} -i\frac{2}{3}\lambda & q_1 & q_2 \\ -q_1^* & \frac{i}{3}\lambda & 0 \\ -q_2^* & 0 & \frac{i}{3}\lambda \end{pmatrix}, \\ V &= U\lambda + \begin{pmatrix} i|q_1|^2 + i|q_2|^2 & iq_{1x} & iq_{2x} \\ iq_{1x}^* & -i|q_1|^2 & -iq_2q_1^* \\ iq_{2x}^* & -iq_2^*q_1 & -i|q_2|^2 \end{pmatrix}.\end{aligned}\quad (\text{A2})$$

The star denotes the complex conjugate. With the trivial seed solutions $q_1[0] = 0, q_2[0] = 0$ and the spectral parameter $\lambda = \lambda_j = a_j + b_j i$ ($j = 1, 2, \dots, N$), the vector eigenfunctions of the linear system Eq. (5) can be written as

$$\Phi_j = \begin{pmatrix} \Phi_{1j} \\ \Phi_{2j} \\ \Phi_{3j} \end{pmatrix} = \begin{pmatrix} e^{-2\theta_j} \\ c_{1j}e^{\theta_j} \\ c_{2j}e^{\theta_j} \end{pmatrix}, \quad \theta_j = \frac{i}{3}\lambda_j x + \frac{i}{3}\lambda_j^2 t, \quad (\text{A3})$$

where c_{1j} and c_{2j} are the coefficients of eigenfunctions, and they are complex parameters. The fundamental bright solitons can be obtained by the following Darboux transformation:

$$\begin{aligned}\Phi[1] &= T[1]\Phi, \quad T[1] = \mathbb{I} - \frac{\lambda_1 - \lambda_1^*}{\lambda - \lambda_1^*} P[1], \\ q_1[1] &= q_1[0] + (\lambda_1^* - \lambda_1)(P[1])_{12}, \\ q_2[1] &= q_2[0] + (\lambda_1^* - \lambda_1)(P[1])_{13}.\end{aligned}\quad (\text{A4})$$

Here, $P[1] = \frac{\Phi_1 \Phi_1^\dagger}{\Phi_1^\dagger \Phi_1}$, where Φ_1 is a special solution at $\lambda = \lambda_1$; a dagger denotes the matrix transpose and complex conjugate. $(P[j])_{ij}$ represents the entry of matrix $P[j]$ in the first row and j column. To obtain a double-hump one soliton, we need to do the second step of the transformation. We employ Φ_2 , which is mapped to $\Phi_2[1] = T[1]|_{\lambda=\lambda_2} \Phi_2$; one double-hump soliton solution can be obtained with the spectral parameter $\lambda_2 = a_1 + ib_2$:

$$\begin{aligned}\Phi[2] &= T[2]\Phi[1], \quad T[2] = \mathbb{I} - \frac{\lambda_2 - \lambda_2^*}{\lambda - \lambda_2^*} P[2], \\ q_1[2] &= q_1[1] + (\lambda_2^* - \lambda_2)(P[2])_{12}, \\ q_2[2] &= q_2[1] + (\lambda_2^* - \lambda_2)(P[2])_{13},\end{aligned}\quad (\text{A5})$$

where $P[2] = \frac{\Phi_2[1]\Phi_2[1]^\dagger}{\Phi_2[1]^\dagger\Phi_2[1]}$. For this case, we choose the coefficients of the eigenfunctions $\Phi_{1,2}$ in Eq. (A3) in the following way: (i) $c_{21} = c_{12} = 0$, and c_{11} and c_{22} are nonzero complex parameters, or (ii) $c_{11} = c_{22} = 0$, and c_{21} and c_{12} are nonzero complex parameters. The corresponding simplified solution

has been present in Eq. (2). Examples of the relevant intensity profiles have been exhibited in Fig. 1.

To study the interaction between nondegenerate solitons, one needs to do a multiple-step transition. For example, by performing the third step of the transition, we employ Φ_3 , which is mapped to $\Phi_3[2] = (T[2]\Phi_3[1])|_{\lambda=\lambda_3}$, with $\Phi_3[1] = (T[1]\Phi_3)|_{\lambda=\lambda_3}$, and then the collision between a double-hump soliton and a single-hump soliton can be obtained with the spectral parameter $\lambda_3 = a_2 + ib_3$:

$$\begin{aligned}\Phi[3] &= T[3]\Phi[2], \quad T[3] = \mathbb{I} - \frac{\lambda_3 - \lambda_3^*}{\lambda - \lambda_3^*} P[3], \\ q_1[3] &= q_1[2] + (\lambda_3^* - \lambda_3)(P[3])_{12}, \\ q_2[3] &= q_2[2] + (\lambda_3^* - \lambda_3)(P[3])_{13},\end{aligned}\quad (\text{A6})$$

where $P[3] = \frac{\Phi_3[2]\Phi_3[2]^\dagger}{\Phi_3[2]^\dagger\Phi_3[2]}$. For this case, we choose the coefficients of the eigenfunctions Φ_3 in Eq. (A3) in the following way: (i) c_{13} and c_{23} are nonzero complex parameters, or (ii) $c_{13} = 0$ or (iii) $c_{23} = 0$, and the coefficients of the eigenfunctions $\Phi_{1,2}$ are the same as those in Eq. (A5). A typical example for this case has been shown in Figs. 3(a1) and 3(a2).

Naturally, by performing the fourth-step of the transformation, one can investigate the interaction between two double-hump solitons. We employ Φ_4 , which is mapped to $\Phi_4[3] = (T[3]\Phi_4[2])|_{\lambda=\lambda_4} = (T[3]T[2]T[1]\Phi_4)|_{\lambda=\lambda_4}$, and then two double-hump solitons' solutions can be obtained as follows with the spectral parameter $\lambda_4 = a_2 + ib_4$:

$$\begin{aligned}q_1[4] &= q_1[3] + (\lambda_4^* - \lambda_4)(P[4])_{12}, \\ q_2[4] &= q_2[3] + (\lambda_4^* - \lambda_4)(P[4])_{13},\end{aligned}\quad (\text{A7})$$

where $P[4] = \frac{\Phi_4[3]\Phi_4[3]^\dagger}{\Phi_4[3]^\dagger\Phi_4[3]}$. For this case, the coefficients of the vector eigenfunctions $\Phi_{3,4}$ are analogous to $\Phi_{1,2}$, i.e., (i) $c_{23} = c_{14} = 0$, and c_{13} and c_{24} are nonzero complex parameters, or (ii) $c_{13} = c_{24} = 0$, and c_{23} and c_{14} are nonzero complex parameters. One typical case has been shown in Figs. 3(b1) and 3(b2).

In general, the N -fold Darboux matrix can be constructed as follows:

$$\begin{aligned}\mathbf{T}_N &= \mathbb{I} - \mathbf{X}_N \mathbf{M}_N^{-1} (\lambda \mathbb{I} - \mathbf{D}_N)^{-1} \mathbf{X}_N^\dagger, \\ \mathbf{X}_N &= [\Phi_1, \Phi_2, \dots, \Phi_N], \\ \mathbf{D}_N &= \text{diag}(\lambda_1^*, \lambda_2^*, \dots, \lambda_N^*), \\ \mathbf{M}_N &= \left(\frac{\Phi_i^\dagger \Phi_j}{\lambda_j - \lambda_i^*} \right)_{1 \leq i, j \leq N},\end{aligned}\quad (\text{A8})$$

and the Bäcklund transformation between old potential functions and new ones can be written as

$$\begin{aligned}q_1[N] &= q_1 + \frac{\det(\mathbf{M}_{N,1})}{\det(\mathbf{M}_N)}, \\ q_2[N] &= q_2 + \frac{\det(\mathbf{M}_{N,2})}{\det(\mathbf{M}_N)},\end{aligned}\quad (\text{A9})$$

where

$$\mathbf{M}_{N,1} = \begin{bmatrix} \mathbf{M}_N & \mathbf{X}_{N,2}^\dagger \\ \mathbf{X}_{N,1} & 0 \end{bmatrix}, \quad \mathbf{M}_{N,2} = \begin{bmatrix} \mathbf{M}_N & \mathbf{X}_{N,3}^\dagger \\ \mathbf{X}_{N,1} & 0 \end{bmatrix}.$$

Here, $\mathbf{X}_{N,i}$ represents the i th row of matrix $\mathbf{X}_N$.

For the two double-hump solitons, we choose the parameters as the following: $\lambda_1 = a_1 + b_1 i$, $\lambda_2 = a_1 + b_2 i$, $c_{21} = 0$, and $c_{12} = 0$, and c_{11} and c_{22} are nonzero complex parameters, which determine the first double-hump soliton; and $\lambda_3 = a_2 + b_3 i$, $\lambda_4 = a_2 + b_4 i$, $c_{23} = 0$, and $c_{14} = 0$, and c_{13} and c_{24} are nonzero complex parameters, which determine the second double-hump soliton. The oscillator for the two solitons is governed by the factors $\sin[(a_1 - a_2)x + (a_1^2 - a_2^2 + b_3^2 - b_1^2)t]$, $\cos[(a_1 - a_2)x + (a_1^2 - a_2^2 + b_3^2 - b_1^2)t]$, $\sin[(a_1 - a_2)x + (a_1^2 - a_2^2 - b_2^2 + b_4^2)t]$, and $\cos[(a_1 - a_2)x + (a_1^2 - a_2^2 - b_2^2 + b_4^2)t]$. The velocity of the solitons is controlled by $x + 2a_j t = \text{const}$, $j = 1$ and 2 , respectively, i.e., the velocity of the soliton is equal to $-2a_j$. Assume that $b_j > 0$ ($j = 1, 2, 3, 4$) and $a_1 > a_2$, and fix the parameters of the first double-hump soliton $x + 2a_1 t = \text{const}$, then $x + 2a_2 t = x + 2a_1 t + 2(a_2 - a_1)t$. If $t \rightarrow +\infty$, we have $x + 2a_2 t \rightarrow -\infty$. Then we see that

$$\Phi_3 \parallel \begin{pmatrix} e^{-i\lambda_j x - i\lambda_j^2 t} \\ c_{13} \\ 0 \end{pmatrix} \rightarrow \begin{pmatrix} 0 \\ c_{13} \\ 0 \end{pmatrix}$$

and

$$\Phi_4 \parallel \begin{pmatrix} e^{-i\lambda_j x - i\lambda_j^2 t} \\ 0 \\ c_{24} \end{pmatrix} \rightarrow \begin{pmatrix} 0 \\ 0 \\ c_{24} \end{pmatrix}.$$

Since the order of iteration for the Darboux transformation can be exchanged, we rewrite

$$T[1] = \mathbb{I} - \frac{\lambda_3 - \lambda_3^*}{\lambda - \lambda_3^*} \frac{\Phi_3 \Phi_3^\dagger}{\Phi_3^\dagger \Phi_3} \rightarrow \text{diag}\left(1, \frac{\lambda - \lambda_3}{\lambda - \lambda_3^*}, 1\right)$$

along the line of $x - a_1 t = \text{const}$ as $t \rightarrow \infty$. It follows that

$$\Phi_4[1] \parallel T[1]|_{\lambda=\lambda_4} \begin{pmatrix} 0 \\ 0 \\ c_{24} \end{pmatrix} \rightarrow \begin{pmatrix} 0 \\ 0 \\ c_{24} \end{pmatrix},$$

which deduces that

$$T[2] \rightarrow \text{diag}\left(1, 1, \frac{\lambda - \lambda_4}{\lambda - \lambda_4^*}\right).$$

Combining the first and second Darboux matrices, we obtain

$$T[2]T[1] \rightarrow \text{diag}\left(1, \frac{\lambda - \lambda_3}{\lambda - \lambda_3^*}, \frac{\lambda - \lambda_4}{\lambda - \lambda_4^*}\right),$$

which yields that

$$\begin{aligned}\Phi_1[2] &= T[2]T[1]|_{\lambda=\lambda_1}\Phi_1 \rightarrow \Phi_1^{[+]} = \begin{pmatrix} e^{-2\theta_1} \\ c_{11}^{[+]}e^{\theta_1} \\ 0 \end{pmatrix}, \\ \Phi_2[2] &= T[2]T[1]|_{\lambda=\lambda_2}\Phi_2 \rightarrow \Phi_2^{[+]} = \begin{pmatrix} e^{-2\theta_2} \\ 0 \\ c_{22}^{[+]}e^{\theta_2} \end{pmatrix},\end{aligned}\tag{A10}$$

where $c_{11}^{[+]} = \frac{\lambda_1 - \lambda_3}{\lambda_1 - \lambda_3^*} c_{11}$ and $c_{22}^{[+]} = \frac{\lambda_2 - \lambda_4}{\lambda_2 - \lambda_4^*} c_{22}$.

Thus, when $t \rightarrow +\infty$, the Darboux matrix tends to

$$\mathbf{T}_4 \rightarrow [\mathbb{I} - \mathbf{X}_2^{[+]}(\mathbf{M}_2^{[+]})^{-1}(\lambda\mathbb{I} - \mathbf{D}_2)^{-1}(\mathbf{X}_2^{[+]})^\dagger] \\ \times \text{diag}\left(1, \frac{\lambda - \lambda_3}{\lambda - \lambda_3^*}, \frac{\lambda - \lambda_4}{\lambda - \lambda_4^*}\right), \quad (\text{A11})$$

where

$$\mathbf{X}_2^{[+]} = [\Phi_1^{[+]}, \Phi_2^{[+]}], \quad \mathbf{D}_2 = \text{diag}(\lambda_1^*, \lambda_2^*),$$

and

$$\mathbf{M}_2^{[+]} = \begin{bmatrix} \frac{1 + |c_{11}^{[+]}|^2 e^{6\text{Re}(\theta_1)}}{\lambda_1 - \lambda_1^*} & \frac{1}{\lambda_2 - \lambda_1^*} \\ \frac{1}{\lambda_1 - \lambda_2^*} & \frac{1 + |c_{22}^{[+]}|^2 e^{6\text{Re}(\theta_2)}}{\lambda_2 - \lambda_2^*} \end{bmatrix}.$$

From the Darboux matrix Eq. (A11), we obtain that the double-hump soliton approaches

$$\begin{aligned} q_1[4] &\rightarrow q_1(x, t; a_1, b_1, b_2, c_{11}^{[+]}, c_{22}^{[+]}) \\ q_2[4] &\rightarrow q_2(x, t; a_1, b_1, b_2, c_{11}^{[+]}, c_{22}^{[+]}) \end{aligned} \quad (\text{A12})$$

as $t \rightarrow \infty$ along the line $x + 2a_1t = \text{const}$, where q_1 and q_2 are given in Eqs. (2). In a similar manner, as $t \rightarrow -\infty$, we have

$$\begin{aligned} q_1[4] &\rightarrow q_1(x, t; a_1, b_1, b_2, c_{11}^{[-]}, c_{22}^{[-]}), \\ q_2[4] &\rightarrow q_2(x, t; a_1, b_1, b_2, c_{11}^{[-]}, c_{22}^{[-]}), \end{aligned} \quad (\text{A13})$$

as $t \rightarrow -\infty$ along the line $x + 2a_1t = \text{const}$, where $c_{11}^{[-]} = \frac{\lambda_1 - \lambda_3^*}{\lambda_1 - \lambda_3} \frac{\lambda_1 - \lambda_4^*}{\lambda_1 - \lambda_4} c_{11}$ and $c_{22}^{[-]} = \frac{\lambda_2 - \lambda_3^*}{\lambda_2 - \lambda_3} \frac{\lambda_2 - \lambda_4^*}{\lambda_2 - \lambda_4} c_{22}$.

Now we consider the asymptotic behavior of the second soliton. Fixing $x + 2a_2t = \text{const}$, as $t \rightarrow \pm\infty$, and we have the asymptotic expressions

$$\begin{aligned} q_1[4] &\rightarrow q_1(x, t; a_2, b_3, b_4, c_{13}^{[\pm]}, c_{24}^{[\pm]}), \\ q_2[4] &\rightarrow q_2(x, t; a_2, b_3, b_4, c_{13}^{[\pm]}, c_{24}^{[\pm]}), \end{aligned} \quad (\text{A14})$$

where $c_{13}^{[+]} = \frac{\lambda_3 - \lambda_1^*}{\lambda_3 - \lambda_1} \frac{\lambda_3 - \lambda_2^*}{\lambda_3 - \lambda_2} c_{13}$, $c_{24}^{[+]} = \frac{\lambda_4 - \lambda_1^*}{\lambda_4 - \lambda_1} \frac{\lambda_4 - \lambda_2^*}{\lambda_4 - \lambda_2} c_{24}$, $c_{13}^{[-]} = \frac{\lambda_3 - \lambda_1}{\lambda_3 - \lambda_1^*} c_{13}$, and $c_{24}^{[-]} = \frac{\lambda_4 - \lambda_2}{\lambda_4 - \lambda_2^*} c_{24}$.

In the general case, the interaction between two-hump solitons is still inelastic. But under the special case,

$$\begin{aligned} |c_{11}^{[+]}| &= |c_{11}^{[-]}| e^{b_1 \delta_1}, & |c_{22}^{[+]}| &= |c_{22}^{[-]}| e^{b_2 \delta_1}, \\ |c_{13}^{[+]}| &= |c_{13}^{[-]}| e^{b_3 \delta_2}, & |c_{24}^{[+]}| &= |c_{24}^{[-]}| e^{b_4 \delta_2}, \end{aligned} \quad (\text{A15})$$

the interaction is elastic. In other words, this is the sufficient condition, Eq. (A15), of elastic interaction of a two-hump soliton.

APPENDIX B: THE EXPRESSIONS OF $\Xi_{1,2}$, $\chi_{1,2}$, AND $\Delta_{1,2,3}$

The general one triple-hump soliton solution in a three-component NLSE is expressed as Eq. (4), where $\Xi_{1,2}$, $\chi_{1,2}$, and $\Delta_{1,2,3}$ are written as

$$\mathbb{E}_1 = \frac{(b_1 - b_2)^2}{(b_1 + b_2)^2} \Delta + |c_1|^2 e^{2\beta_1} + |c_2|^2 e^{2\beta_2} + |c_3|^2 e^{2\beta_3} + |c_4|^2 e^{2\beta_4} + |c_5|^2 e^{2\beta_5} + |c_6|^2 e^{2\beta_6} + |c_7|^2 e^{2\beta_7} + |c_8|^2 e^{2\beta_8} + |c_9|^2 e^{2\beta_9} + |c_{10}|^2 e^{2\beta_{10}} + |c_{11}|^2 e^{2\beta_{11}} + |c_{12}|^2 e^{2\beta_{12}} + |c_{13}|^2 e^{2\beta_{13}} + |c_{14}|^2 e^{2\beta_{14}} + |c_{15}|^2 e^{2\beta_{15}} + |c_{16}|^2 e^{2\beta_{16}} + |c_{17}|^2 e^{2\beta_{17}} + |c_{18}|^2 e^{2\beta_{18}} + |c_{19}|^2 e^{2\beta_{19}} + |c_{20}|^2 e^{2\beta_{20}} + |c_{21}|^2 e^{2\beta_{21}} + |c_{22}|^2 e^{2\beta_{22}} + |c_{23}|^2 e^{2\beta_{23}} + |c_{24}|^2 e^{2\beta_{24}} + |c_{25}|^2 e^{2\beta_{25}} + |c_{26}|^2 e^{2\beta_{26}} + |c_{27}|^2 e^{2\beta_{27}} + |c_{28}|^2 e^{2\beta_{28}} + |c_{29}|^2 e^{2\beta_{29}} + |c_{30}|^2 e^{2\beta_{30}} + |c_{31}|^2 e^{2\beta_{31}} + |c_{32}|^2 e^{2\beta_{32}} + |c_{33}|^2 e^{2\beta_{33}} + |c_{34}|^2 e^{2\beta_{34}} + |c_{35}|^2 e^{2\beta_{35}} + |c_{36}|^2 e^{2\beta_{36}} + |c_{37}|^2 e^{2\beta_{37}} + |c_{38}|^2 e^{2\beta_{38}} + |c_{39}|^2 e^{2\beta_{39}} + |c_{40}|^2 e^{2\beta_{40}} + |c_{41}|^2 e^{2\beta_{41}} + |c_{42}|^2 e^{2\beta_{42}} + |c_{43}|^2 e^{2\beta_{43}} + |c_{44}|^2 e^{2\beta_{44}} + |c_{45}|^2 e^{2\beta_{45}} + |c_{46}|^2 e^{2\beta_{46}} + |c_{47}|^2 e^{2\beta_{47}} + |c_{48}|^2 e^{2\beta_{48}} + |c_{49}|^2 e^{2\beta_{49}} + |c_{50}|^2 e^{2\beta_{50}} + |c_{51}|^2 e^{2\beta_{51}} + |c_{52}|^2 e^{2\beta_{52}} + |c_{53}|^2 e^{2\beta_{53}} + |c_{54}|^2 e^{2\beta_{54}} + |c_{55}|^2 e^{2\beta_{55}} + |c_{56}|^2 e^{2\beta_{56}} + |c_{57}|^2 e^{2\beta_{57}} + |c_{58}|^2 e^{2\beta_{58}} + |c_{59}|^2 e^{2\beta_{59}} + |c_{60}|^2 e^{2\beta_{60}} + |c_{61}|^2 e^{2\beta_{61}} + |c_{62}|^2 e^{2\beta_{62}} + |c_{63}|^2 e^{2\beta_{63}} + |c_{64}|^2 e^{2\beta_{64}} + |c_{65}|^2 e^{2\beta_{65}} + |c_{66}|^2 e^{2\beta_{66}} + |c_{67}|^2 e^{2\beta_{67}} + |c_{68}|^2 e^{2\beta_{68}} + |c_{69}|^2 e^{2\beta_{69}} + |c_{70}|^2 e^{2\beta_{70}} + |c_{71}|^2 e^{2\beta_{71}} + |c_{72}|^2 e^{2\beta_{72}} + |c_{73}|^2 e^{2\beta_{73}} + |c_{74}|^2 e^{2\beta_{74}} + |c_{75}|^2 e^{2\beta_{75}} + |c_{76}|^2 e^{2\beta_{76}} + |c_{77}|^2 e^{2\beta_{77}} + |c_{78}|^2 e^{2\beta_{78}} + |c_{79}|^2 e^{2\beta_{79}} + |c_{80}|^2 e^{2\beta_{80}} + |c_{81}|^2 e^{2\beta_{81}} + |c_{82}|^2 e^{2\beta_{82}} + |c_{83}|^2 e^{2\beta_{83}} + |c_{84}|^2 e^{2\beta_{84}} + |c_{85}|^2 e^{2\beta_{85}} + |c_{86}|^2 e^{2\beta_{86}} + |c_{87}|^2 e^{2\beta_{87}} + |c_{88}|^2 e^{2\beta_{88}} + |c_{89}|^2 e^{2\beta_{89}} + |c_{90}|^2 e^{2\beta_{90}} + |c_{91}|^2 e^{2\beta_{91}} + |c_{92}|^2 e^{2\beta_{92}} + |c_{93}|^2 e^{2\beta_{93}} + |c_{94}|^2 e^{2\beta_{94}} + |c_{95}|^2 e^{2\beta_{95}} + |c_{96}|^2 e^{2\beta_{96}} + |c_{97}|^2 e^{2\beta_{97}} + |c_{98}|^2 e^{2\beta_{98}} + |c_{99}|^2 e^{2\beta_{99}} + |c_{100}|^2 e^{2\beta_{100}} + |c_{101}|^2 e^{2\beta_{101}} + |c_{102}|^2 e^{2\beta_{102}} + |c_{103}|^2 e^{2\beta_{103}} + |c_{104}|^2 e^{2\beta_{104}} + |c_{105}|^2 e^{2\beta_{105}} + |c_{106}|^2 e^{2\beta_{106}} + |c_{107}|^2 e^{2\beta_{107}} + |c_{108}|^2 e^{2\beta_{108}} + |c_{109}|^2 e^{2\beta_{109}} + |c_{110}|^2 e^{2\beta_{110}} + |c_{111}|^2 e^{2\beta_{111}} + |c_{112}|^2 e^{2\beta_{112}} + |c_{113}|^2 e^{2\beta_{113}} + |c_{114}|^2 e^{2\beta_{114}} + |c_{115}|^2 e^{2\beta_{115}} + |c_{116}|^2 e^{2\beta_{116}} + |c_{117}|^2 e^{2\beta_{117}} + |c_{118}|^2 e^{2\beta_{118}} + |c_{119}|^2 e^{2\beta_{119}} + |c_{120}|^2 e^{2\beta_{120}} + |c_{121}|^2 e^{2\beta_{121}} + |c_{122}|^2 e^{2\beta_{122}} + |c_{123}|^2 e^{2\beta_{123}} + |c_{124}|^2 e^{2\beta_{124}} + |c_{125}|^2 e^{2\beta_{125}} + |c_{126}|^2 e^{2\beta_{126}} + |c_{127}|^2 e^{2\beta_{127}} + |c_{128}|^2 e^{2\beta_{128}} + |c_{129}|^2 e^{2\beta_{129}} + |c_{130}|^2 e^{2\beta_{130}} + |c_{131}|^2 e^{2\beta_{131}} + |c_{132}|^2 e^{2\beta_{132}} + |c_{133}|^2 e^{2\beta_{133}} + |c_{134}|^2 e^{2\beta_{134}} + |c_{135}|^2 e^{2\beta_{135}} + |c_{136}|^2 e^{2\beta_{136}} + |c_{137}|^2 e^{2\beta_{137}} + |c_{138}|^2 e^{2\beta_{138}} + |c_{139}|^2 e^{2\beta_{139}} + |c_{140}|^2 e^{2\beta_{140}} + |c_{141}|^2 e^{2\beta_{141}} + |c_{142}|^2 e^{2\beta_{142}} + |c_{143}|^2 e^{2\beta_{143}} + |c_{144}|^2 e^{2\beta_{144}} + |c_{145}|^2 e^{2\beta_{145}} + |c_{146}|^2 e^{2\beta_{146}} + |c_{147}|^2 e^{2\beta_{147}} + |c_{148}|^2 e^{2\beta_{148}} + |c_{149}|^2 e^{2\beta_{149}} + |c_{150}|^2 e^{2\beta_{150}} + |c_{151}|^2 e^{2\beta_{151}} + |c_{152}|^2 e^{2\beta_{152}} + |c_{153}|^2 e^{2\beta_{153}} + |c_{154}|^2 e^{2\beta_{154}} + |c_{155}|^2 e^{2\beta_{155}} + |c_{156}|^2 e^{2\beta_{156}} + |c_{157}|^2 e^{2\beta_{157}} + |c_{158}|^2 e^{2\beta_{158}} + |c_{159}|^2 e^{2\beta_{159}} + |c_{160}|^2 e^{2\beta_{160}} + |c_{161}|^2 e^{2\beta_{161}} + |c_{162}|^2 e^{2\beta_{162}} + |c_{163}|^2 e^{2\beta_{163}} + |c_{164}|^2 e^{2\beta_{164}} + |c_{165}|^2 e^{2\beta_{165}} + |c_{166}|^2 e^{2\beta_{166}} + |c_{167}|^2 e^{2\beta_{167}} + |c_{168}|^2 e^{2\beta_{168}} + |c_{169}|^2 e^{2\beta_{169}} + |c_{170}|^2 e^{2\beta_{170}} + |c_{171}|^2 e^{2\beta_{171}} + |c_{172}|^2 e^{2\beta_{172}} + |c_{173}|^2 e^{2\beta_{173}} + |c_{174}|^2 e^{2\beta_{174}} + |c_{175}|^2 e^{2\beta_{175}} + |c_{176}|^2 e^{2\beta_{176}} + |c_{177}|^2 e^{2\beta_{177}} + |c_{178}|^2 e^{2\beta_{178}} + |c_{179}|^2 e^{2\beta_{179}} + |c_{180}|^2 e^{2\beta_{180}} + |c_{181}|^2 e^{2\beta_{181}} + |c_{182}|^2 e^{2\beta_{182}} + |c_{183}|^2 e^{2\beta_{183}} + |c_{184}|^2 e^{2\beta_{184}} + |c_{185}|^2 e^{2\beta_{185}} + |c_{186}|^2 e^{2\beta_{186}} + |c_{187}|^2 e^{2\beta_{187}} + |c_{188}|^2 e^{2\beta_{188}} + |c_{189}|^2 e^{2\beta_{189}} + |c_{190}|^2 e^{2\beta_{190}} + |c_{191}|^2 e^{2\beta_{191}} + |c_{192}|^2 e^{2\beta_{192}} + |c_{193}|^2 e^{2\beta_{193}} + |c_{194}|^2 e^{2\beta_{194}} + |c_{195}|^2 e^{2\beta_{195}} + |c_{196}|^2 e^{2\beta_{196}} + |c_{197}|^2 e^{2\beta_{197}} + |c_{198}|^2 e^{2\beta_{198}} + |c_{199}|^2 e^{2\beta_{199}} + |c_{200}|^2 e^{2\beta_{200}} + |c_{201}|^2 e^{2\beta_{201}} + |c_{202}|^2 e^{2\beta_{202}} + |c_{203}|^2 e^{2\beta_{203}} + |c_{204}|^2 e^{2\beta_{204}} + |c_{205}|^2 e^{2\beta_{205}} + |c_{206}|^2 e^{2\beta_{206}} + |c_{207}|^2 e^{2\beta_{207}} + |c_{208}|^2 e^{2\beta_{208}} + |c_{209}|^2 e^{2\beta_{209}} + |c_{210}|^2 e^{2\beta_{210}} + |c_{211}|^2 e^{2\beta_{211}} + |c_{212}|^2 e^{2\beta_{212}} + |c_{213}|^2 e^{2\beta_{213}} + |c_{214}|^2 e^{2\beta_{214}} + |$$

$$\begin{aligned}
\Delta_1 &= (b_2 - b_3) \left[\frac{(b_1 - b_2)^2 (b_1 - b_3)}{(b_1 + b_2)^2 (b_1 + b_3)} - |c_1|^2 e^{2\beta_1} \right] \\
&\quad + (b_2 + b_3) \left[\frac{b_3 - b_1}{b_3 + b_1} + |c_1|^2 e^{2\beta_1} \right] |c_2|^2 e^{2\beta_2}, \\
\Delta_2 &= \frac{(b_1 - b_2)(b_2 - b_3)}{(b_1 + b_2)(b_1 + b_3)} - \frac{b_2 + b_3}{b_1 + b_3} |c_2|^2 e^{2\beta_2}, \\
\Delta_3 &= \frac{(b_2^2 - b_1^2)(b_1 - b_3)}{(b_1 + b_2)^2 (b_1)}
\end{aligned}$$

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