Coherently controlled entanglement generation in a binary Bose-Einstein condensate

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Abstract – Considering a two-component Bose-Einstein condensate in a double-well potential, a method to generate a Bell state consisting of two spatially separated condensates is suggested. For repulsive interactions, the required tunnelling control is achieved numerically by varying the amplitude of a sinusoidal potential difference between the wells. Both numerical and analytical calculations reveal the emergence of a highly entangled mesoscopic state.

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In a recent ground-breaking experiment self-trapping was observed for small Bose-Einstein condensates (BECs) in a double-well potential \cite{1}. A major experimental goal in this system is the generation of mesoscopic entangled number states like “Schrödinger-cat” states \cite{1}. Unfortunately, in a double well the ground state of a condensate with repulsive interaction is not the desired “Schrödinger-cat” state (e.g. ref. \cite{2}). On the other hand, using attractive condensates leads to the problem of instability (e.g. ref. \cite{3}). However, a two-component repulsive condensate in a double well is investigated in this letter, because it offers the possibility to increase the entanglement of the ground state drastically.

To control the entanglement generation, we employ a method recently suggested to induce the phase transition between a superfluid and a Mott insulator \cite{4} by applying time-periodic potential differences \cite{5,6}. While on the single particle level the effects predicted for periodically driven systems \cite{7,8} have not yet been realised experimentally, the coherent control by changing the amplitude of periodic force fields might be easier to achieve for BECs rather than for single particles.

Suggestions to produce mesoscopic entangled states, involve coherent scattering of far detuned light fields \cite{9}; several groups suggest to produce mesoscopic entanglement with condensates by manipulating the interaction between the particles, or by controlling the dynamics of the system \cite{10–15}. While the experimental realisation of such mesoscopic entangled states is still a challenge of current fundamental research, possible applications include quantum computing or quantum information processing.

In this letter we study a double-well potential similar to the experiment \cite{1}, but with a two-component condensate (for various aspects investigated for two-component BECs see, e.g., refs. \cite{16–19}). To motivate the mechanism by which we control the entanglement generation, we start with the two-particle case, because it exhibits all crucial features, before we extend both the numerical and analytical calculations to $N$ particles.

The model. – We study a symmetric double-well potential filled with Bose particles at very low temperature. Adopting the common two-mode approximation, assuming only on-site interactions and denoting the tunnelling splitting between the two lowest single particle energy states with $\hbar \Omega$, we use a well-established model describing the dynamics of a one-species BEC in double-well potential \cite{20}. This model gives a good qualitative description of recent experiments \cite{1}. Extending the model to a binary condensate consisting of $N$ particles of species A and $N$ particles of species B leads to the
Hamiltonian

\[
\hat{H} = -\frac{\hbar \Omega}{2}(\hat{a}_1^\dagger \hat{a}_1 + \hat{a}_2^\dagger \hat{a}_2) + \hbar \kappa_A (\hat{a}_1^\dagger \hat{a}_1 \hat{a}_1 + \hat{a}_1^\dagger \hat{a}_2 \hat{a}_2) \\
- \frac{\hbar \Omega}{2}(\hat{b}_1^\dagger \hat{b}_1 + \hat{b}_2^\dagger \hat{b}_2) + \hbar \kappa_B (\hat{b}_1^\dagger \hat{b}_1 \hat{b}_1 + \hat{b}_2^\dagger \hat{b}_2 \hat{b}_2) \\
+ 2\hbar \kappa_{AB} (\hat{a}_1^\dagger \hat{b}_1 \hat{b}_1^\dagger \hat{a}_2 + \hat{a}_1^\dagger \hat{b}_2 \hat{b}_2^\dagger \hat{a}_1) \\
+ \hbar \mu f(t) (\hat{a}_1^\dagger \hat{a}_1 - \hat{a}_2^\dagger \hat{a}_2 + \hat{b}_1 - \hat{b}_2),
\]

(1)

where \(\hat{a}_i^{(t)}\) and \(\hat{b}_i^{(t)}\) are the annihilation (creation) operators for a boson of species A or B in the \(i\)-th well \((i = 1, 2)\) and \(\hbar \mu f(t)\) specifies an externally applied potential difference between the two wells. The relation between the interaction parameters \(\kappa_A, \kappa_B\) and \(\kappa_{AB}\) can be found, e.g., in ref. [2]. We use the Fock-basis \(|i, N - i\rangle_A|j, N - j\rangle_B\) \((i, j = \{0, 1\})\), where \(i/j\) are the numbers of A/B particles occupying the first well and \((N - i)/(N - j)\) the number of A/B particles occupying the second well.

In these notations, the Bell state reads:

\[
\psi_{\text{Bell}} = \frac{1}{\sqrt{2}}(|N, 0\rangle + |0, N\rangle) \\
\equiv \frac{1}{\sqrt{2}}(|N, 0\rangle_A|0, N\rangle_B + |0, N\rangle_A|N, 0\rangle_B).
\]

(2)

In this state, the two condensates are maximally entangled with each other in the sense that each condensate is in a superposition of being in both wells and if one condensate is measured in one well, the other condensate is measured in the other well.

Similar spatially entangled states are known from the EPR paradox and are of special interest in the field of quantum cryptography, where especially entangled photons are used [21,22]. Looking at the \(N\)-particle dynamics, we have to take \((N + 1)^2\) states into account. If we exemplarily study the case of only two distinguishable particles A and B without driving the well \((\mu = 0)\), the system can be solved analytically. While we assume short range interaction between the two particles typical for cold atoms, the resulting entanglement is similar to the case of two electrons in a double quantum dot [23]. Here, we use the basis \(\{1, 1, 1, 0, 0, 1, 0, 0\}\) to obtain the system’s Hamiltonian as a \(4 \times 4\) matrix,

\[
\hat{H}_2 = \hbar \Omega \begin{pmatrix}
  u & -1/2 & -1/2 & 0 \\
  -1/2 & 0 & 0 & -1/2 \\
  -1/2 & 0 & 0 & -1/2 \\
  0 & -1/2 & -1/2 & u
\end{pmatrix},
\]

(3)

where \(u = \kappa_{AB}/\Omega\) is the interaction parameter. When repulsive interaction is assumed \((u > 0)\), the ground state and first excited state of the system are given by

\[
y_0 = \mathcal{N} \begin{pmatrix}
  1 \\
  \frac{-1 + \sqrt{u^2 + 4}}{2} \\
  \frac{-1 - \sqrt{u^2 + 4}}{2} \\
  1
\end{pmatrix}, \quad y_1 = \frac{1}{\sqrt{2}} \begin{pmatrix}
  0 \\
  -1 \\
  1 \\
  0
\end{pmatrix}
\]

(4)

with norm \(\mathcal{N} = (4 + u^2 + u\sqrt{u^2 + 4})^{-1/2}\) and corresponding energies \(E_0 = \hbar \Omega (u - \sqrt{u^2 + 4}/2)\) and \(E_1 = 0\). In the case of a very large interaction parameter the ground state approaches the Bell state \((\text{eq. } (2)\text{ with } N = 1)\),

\[
y_0 \xrightarrow{u \to \infty} \frac{1}{\sqrt{2}} (1 0 1 0)^\dagger = \frac{1}{\sqrt{2}} ((0, 1) + (1, 0)),
\]

(5)

and becomes a maximally entangled state. On the other hand with no interaction \((u = 0)\) the ground state is given by \(1/2 \cdot (1 \ 1 \ 1 \ 1)^\dagger = 1/2 \cdot ((1, 1) + (0, 1) + (1, 0) + (0, 0))\), which can be written as a product state \(y_0(u = 0)) = 1/2 ((1, 0)_A + (0, 1)_A) \cdot ((1, 0)_B + (0, 1)_B)\) and therefore is not entangled at all. Furthermore, for very large interaction parameter \(u\) the ground state is quasi-degenerate. This is also true for the \(N\)-particle system with the states \(|(N, 0) \pm |0, N\rangle\rangle/\sqrt{2}\). Because of this quasi-degeneracy one cannot generate a Bell state by just generating a binary-BEC in a very deep double-well potential; the system could be in any superposition of those two states. If, however, one starts in the ground state of a not so deep well, the Hamiltonian does not allow transitions between symmetric and antisymmetric states (cf. ref. [24]). Thus, thermal excitations to the (antisymmetric) first excited state can be discarded.

The double-well potential is now modulated periodically in time by the force \(f\) indicated in \(\text{eq. } (1)\). As the new Hamiltonian is periodic in time, Floquet theory can be applied. As shown in refs. [25,26], for sufficiently high frequencies \((\omega/\Omega \gg 1)\) the system behaves like an undriven system with a rescaled effective tunnelling frequency \(\epsilon_{\text{eff}} = \Omega \cdot J_0(2\mu/\omega)\) or equivalently with the effective interaction parameter

\[
u \to u_{\text{eff}} = \frac{u}{J_0 \left(\frac{2\mu}{\omega}\right)},
\]

(6)

where \(J_0\) denotes the ordinary Bessel function of order zero. The system is coherently controlled by driving it with linearly in time increasing parameter \(2\mu/\omega\) (in this case \(2\mu/\omega \simeq 2.405\cdot \tau/\tau_{\text{max}}\), where \(\tau = \Omega \cdot t\) is a dimensionless time variable). At the time \(\tau_{\text{max}}\) the parameter reaches the first zero of \(J_0\); the effective interaction becomes very large. Similar effects could be achieved without driving by adiabatically increasing the depth of the wells. Compared with this coherent control mechanism, the feed-back loops of optimal control theory [27] would allow to further optimise entanglement generation for given experimental parameters.

To quantify the entanglement of a state \(|\Psi\rangle\) we use the fidelity, i.e., the probability of the system to be in the Bell state given by the square of the state’s overlap with \(\psi_{\text{Bell}}\):

\[
P_{\text{Bell}} = |\langle \psi_{\text{Bell}} | \Psi \rangle|^2.
\]

(7)

Because our goal are fidelities close to one, more sophisticated entanglement measures (see, e.g., refs. [22,28]), which identify much less entangled states, are not
Fig. 1: Fidelity (eq. (7), probability of the system to be in the Bell state) for the system with $N = 1$ particle of each species. After initialising it in the ground state the driving parameter $2\mu/\omega$ was linearly increased in the time $\tau$ until reaching the first zero of $J_0(2\mu/\omega)$ with $(2\mu/\omega)_{\text{max}} \approx 2.405$; the interaction parameter was $\kappa = \kappa_{\text{AB}}/\Omega = 0.5$. The driving frequency used is $\omega/\Omega = 10$ (upper curve) and $\omega/\Omega = 1.2$ (lower curve). The increase with the higher frequency is in almost perfect agreement with the ideal curve indicated by the dashed line. For the lower frequency the increase does not work and the Bell state is not reached.

Analytical calculation for the Bell state. – In the spirit of refs. [2,10] we use the mean-field equations, in order to extend the two-particle dynamics to many bosons, and construct a symmetrised superposition of two mean-field states. Within the mean-field approximation the system is reduced to four amplitudes describing the dynamics of condensate $A$ ($B$), such that $|a_1|^2 = 1 - |a_2|^2$ ($|b_1|^2 = 1 - |b_2|^2$) is the expected fraction of condensate $A$ ($B$) in the first well. The equations of motion for two condensates with equal number of particles ($N_A = N_B = N$) are given by

$$i\dot{a}_1 = -\frac{1}{2}a_2 + 2N\frac{\kappa_A}{\Omega}|a_1|^2a_1 + 4N\frac{\kappa_{\text{AB}}}{\Omega}|b_1|^2a_1,$$

$$i\dot{a}_2 = -\frac{1}{2}a_1 + 2N\frac{\kappa_A}{\Omega}|a_2|^2a_2 + 4N\frac{\kappa_{\text{AB}}}{\Omega}|b_2|^2a_2,$$

$$i\dot{b}_1 = -\frac{1}{2}b_2 + 2N\frac{\kappa_B}{\Omega}|b_1|^2b_1 + 4N\frac{\kappa_{\text{AB}}}{\Omega}|a_1|^2b_1,$$

$$i\dot{b}_2 = -\frac{1}{2}b_1 + 2N\frac{\kappa_B}{\Omega}|b_2|^2b_2 + 4N\frac{\kappa_{\text{AB}}}{\Omega}|a_2|^2b_2,$$

where the dot means the derivative with respect to $\tau = \Omega \cdot t$ and $\kappa_A$, $\kappa_B$ and $\kappa_{\text{AB}}$ are the interaction parameters given in eq. (1). In the following the intra-condensate interactions are equalised and the abbreviations $\alpha \equiv 2N\kappa_A/N\Omega$ and $\beta \equiv 2N\kappa_{\text{AB}}/N\Omega$ are used. Experimentally, a similar situation could be achieved by choosing the numbers of particles in the condensates, such that $N_A\kappa^2 = N_B\kappa^2$. The inter-condensate interaction parameter $\beta$ can be modulated by a Feshbach-resonance [29].

To find stationary state solutions we apply the ansatz

$$a_{1,2} = \psi_{1,2} \exp(-i\mu t) \quad \text{and} \quad b_{1,2} = \phi_{1,2} \exp(-i\mu t).$$

One finds two kinds of solutions, a balanced one, where the condensates are equally distributed on the two wells ($\psi_1^2 = \psi_2^2 = \phi_1^2 = \phi_2^2 = 1/2$) and an unbalanced solution, where condensate $A/B$ is mainly in the first/second well. For $(2\beta - \alpha) > 0$ the amplitudes in this case are

$$\psi_{1,2}/\beta_2/1 = \frac{1}{\sqrt{2}} \left( 1 \pm \sqrt{1 - \frac{1}{(2\beta - \alpha)^2}} \right)^{1/2}.$$  

A second unbalanced solution is obtained by interchanging the indices 1 and 2. If $(2\beta - \alpha)$ is negative, $\psi_1$ and $\psi_2$ in eq. (13) have opposite signs. As the resulting state is not the ground state in the limit of large interactions necessary for entanglement generation, this solution can be discarded here. The mean-field Hamiltonian corresponding to eqs. (8)-(11) reads:

$$H = \frac{\hbar \Omega}{2} \left[ -(a_1a_2^* + a_2a_1^*) + \alpha(|a_1|^4 + |a_2|^4) \right] + \frac{\hbar \Omega}{2} \left[ -(b_1b_2^* + b_2b_1^*) + \alpha(|b_1|^4 + |b_2|^4) \right] + \hbar \Omega \alpha (|a_1|^2|b_1|^2 + |a_2|^2|b_2|^2).$$

The energies for the balanced and unbalanced states per particle therefore are given by

$$\frac{E_{\text{bal}}}{N\Omega} = \frac{\alpha + \beta}{2} \pm 1, \quad \frac{E_{\text{unbal}}}{N\Omega} = \alpha + \frac{\alpha - 3\beta}{2(2\beta - \alpha)^2}. \quad (15)$$

When is the unbalanced state the ground state? It is, if $E_{\text{unbal}} < E_{\text{bal}}$, leading to

$$\alpha - \beta + 2 + \frac{\alpha - 3\beta}{(2\beta - \alpha)^2} < 0. \quad (16)$$

When again by driving the double well the effective interactions are strongly increased, see eq. (6), only the difference of the intra- and inter-condensate interactions $\alpha$ and $\beta$ determine whether the balanced or the unbalanced state is the ground state.

To compare the nonlinear mean-field dynamics to the $N$-particle dynamics, we need the $N$-particle counterparts

\footnote{One possible realisation of such a binary condensate could be two hyperfine-states of $^{87}$Rb, for which Feshbach-resonances between different hyperfine-states have been observed experimentally [29].}
of the mean-field states. By using the atomic coherent states [2]

\[ |\vartheta, \varphi\rangle_A = \sum_{n_A=0}^{N} \binom{N}{n_A}^{1/2} \cos^{n_A}(\vartheta/2) \sin^{N-n_A}(\vartheta/2) \]

\[ \times e^{i(N-n_A)\varphi}|n_A, N-n_A\rangle_A \]

(17)

for condensate A (and B analogously) with \( \cos(\vartheta/2) = \psi_1 \), \( \sin(\vartheta/2) = \psi_2 \) and phase \( \varphi = 0 \) we can construct a symmetrised superposition of two unbalanced states

\[ |\Psi\rangle = N \left( |\vartheta, 0\rangle_A |\vartheta, 0\rangle_B + |\vartheta, 0\rangle_A |\vartheta, 0\rangle_B \right) \]

(18)

where the overlined term is obtained by replacing \( |n, N-n\rangle_{A/B} \) by \( |N-n, n\rangle_{A/B} \). The probability \( P_{\text{Bell}} \) is then given by

\[ P_{\text{Bell}} = 2N^2 \left( \psi_1^{2N} + \psi_2^{2N} \right)^2. \]

(19)

For large \( N \) the above expression is approximated by

\[ P_{\text{Bell}} \approx \exp \left( \frac{-N J_0 (2\mu)^2}{2(2\beta - \alpha)^2} \right). \]

(20)

**N-particle dynamics and entanglement generation.** – To demonstrate entanglement generation numerically, we use 100 particles (50 of each kind), which is close to the experimentally reachable domain [30]. In experiments the decision for small condensates makes sense, because decoherence times induced by particle loss are long. Again we prepare the system in the ground state and drive it with linearly in time \( \tau \) increasing driving parameter \( 2\mu/\omega \) up to the first zero of \( J_0 \). Figure 2 shows the emergence of the highly entangled state under the influence of a periodic force field. The parameters are chosen depending on the desired behaviour: generation of entanglement on very short time scales (fig. 2, top), good agreement with the analytical calculation (eq. (20)) or generation of entanglement starting with a balanced initial state (see fig. 3). Although the mean-field

![Fig. 2: Fidelity (7) for the system with \( N = 50 \) particles of each species (cf. fig. 1). After initialising it in the ground state the driving parameter \( 2\mu/\omega \) was linearly increased until reaching the first zero of \( J_0(2\mu/\omega) \); the used driving frequency was \( \omega/\Omega = 40 \). The upper plot shows that generating the Bell state is possible on a short time scale with a fidelity higher than 96%; the interaction parameters were \( \kappa_A/\Omega = \kappa_B/\Omega = 0.1 \) and \( \kappa_{AB}/\Omega = 0.115 \). Inset: if the interaction parameters are chosen such that the initial state is not bimodal (see fig. 3) with \( \kappa_A/\Omega = \kappa_B/\Omega = 0.1 \) and \( \kappa_{AB}/\Omega = 0.11 \), it takes longer to adiabatically reach the Bell state. The fidelity is again very high (>97.5%). Lower panel: although the unbalanced state is already quite entangled when it becomes the ground state for \( \tau \approx 80 \) (cf. eq. (16)), the analytic curve indicated by the dashed line (cf. eq. (20)) gives a good qualitative description of this process. The used interaction parameters here are \( \kappa_A/\Omega = \kappa_B/\Omega = 0.0016 \) and \( \kappa_{AB}/\Omega = 0.0133 \); the fidelity is >97.5%.

![Fig. 3: Probabilities of the system to be in the Fock state \( |i, j\rangle \), where \( i \) and \( j \) are the Fock indices. The system was initialised in the ground state, which is in this case (same parameters as in the inset of fig. 2) a quite balanced state (upper panel: system at \( \tau = 0 \)), and then driven with linearly increasing driving parameter \( 2\mu/\omega \). With increasing time the system’s state becomes more and more unbalanced (lower panel: system at \( \tau = 250 \)) until it reaches the Bell state at \( \tau_{\text{max}} = 500 \) (cf. inset of fig. 2).
equations are valid in the limit \(N \to \infty\) and \(\kappa \to 0\) such that \(N\kappa\) is constant, the predictions obtained by our calculations qualitatively agree with the numerical results as long as \((2\beta - \alpha)\) is not too high in the regime, where the balanced state is the ground state.

In fig. 3 we show the occupation of states by plotting the Fock occupation numbers \(i\) and \(j\) and the corresponding probabilities \(|\langle \Psi|i,j\rangle|^2\). The probabilities move to the outer (more unbalanced) states while increasing the parameter \(2\mu/\omega\) and at the end only the two states \(|N,0\rangle\) and \(|0,N\rangle\) are nearly equally occupied. So a quite balanced state is changed into a totally unbalanced and maximally entangled state.

Within the validity of the two-mode approximation (cf. ref. [15]), the fidelity \(7)\) can be raised by using higher driving frequencies \(\omega/\Omega\). As the entanglement generation involves adiabatic following of Floquet states (see ref. [5] and references therein), the time-scales on which the amplitude will have to be changed in future experiments can be deduced from calculating the quasienergies (fig. 4). For higher particle numbers than chosen in fig. 4 and not too high driving frequencies, avoided crossings occur which are not important for entanglement generation as long as the driving frequency still is in the high frequency limit (fig. 2; see also ref. [5]). To find the best parameters for a given experiment with respect to low decoherence and large entanglement and to optimise the protocol \(\mu(\tau)\), i.e., the way the driving amplitude is changed in time, one might employ the methods of optimal control theory [27].

**Conclusion.** – We demonstrated that the generation of a highly entangled state with a binary condensate in a double-well potential is possible. By driving the double-well with an increasing amplitude the ground state is turned into a *Bell state* provided that the inter-condensate interaction is first slightly increased via a Feshbach-resonance. The fidelities of our results within the model of eq. (1) are even on short time scales above 96% which is considerably higher than the fidelity obtained in similar models for single-species condensates [13–15].

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