Strong parity effect of particle number in the interference fringes of Bose-Einstein condensates released from a double-well potential

Hongwei Xiong\textsuperscript{1,2} and Shujuan Liu\textsuperscript{1,2}

\textsuperscript{1}State Key Laboratory of Magnetic Resonance and Atomic and Molecular Physics, Wuhan Institute of Physics and Mathematics, Chinese Academy of Sciences, Wuhan 430071, P. R. China
\textsuperscript{2}Center for Cold Atom Physics, Chinese Academy of Sciences, Wuhan 430071, P. R. China

(Dated: July 7, 2009)

We study the parity effect of the particle number in the interference fringes of a Bose-Einstein condensate released from a double-well potential. For a coherently splitting condensate in the double-well potential, with a decoupled two-mode Bose-Hubbard model, there is well-known phase diffusion because of interatomic interactions. After a specific holding time of the double-well potential, the phase diffusion will make the interference patterns in the density distribution depend strongly on the parity of the total particle number by further overlapping two condensates. This parity effect originates from the quantized relative phase about the total particle number. The experimental scheme to observe this “even-odd” effect of the particle number is discussed.

I. INTRODUCTION

For most classical systems with a large number of particles, adding or decreasing one particle will not change the fundamental properties of the system. For the quantum system of dilute Bose or Fermi gases in a harmonic trap with a large number of particles, it is also widely believed that adding or decreasing one particle will not change the thermodynamic and dynamic properties of the system. For example, for a Bose-Einstein condensate (BEC) described by the Gross-Pitaevskii equation within the mean-field model, adding or decreasing one particle plays a negligible role in the coupling parameter (which is proportional to the overall particle number) for the order parameter. Recently, it is found theoretically that for specific cold atomic system beyond the mean-field model, there is a parity effect of the particle number which means that adding or decreasing one particle would have different quantum effect. The parity effect has been studied theoretically for the quantum decay of Josephson $\pi$ states \cite{1}, the Berry phase of two-species Bose-Einstein condensates (BECs) \cite{2}, the tunnel splitting in the Josephson model \cite{3}, the macroscopic superpositions of phase states \cite{4}, and more recently in the mesoscopic quantum switching \cite{5}.

To our best knowledge, the parity effect of particle number has not been observed experimentally for cold atoms, although the parity-dependent tunneling splitting was observed for magnetic molecular clusters \cite{6,7,8,9}. It is still an open question to observe the parity-dependent effect for cold atomic system. For example, it is very difficult for an experiment to observe the parity-dependent effect in the tunneling splitting, because the tunneling splitting is extremely small \cite{5}.

The quantum interference between two condensates in a double-well potential \cite{10,11,12,13,14,15,16,17,18,19,20,21,22} has renewed interest in the experimental and theoretical studies of the macroscopic quantum coherence effect. In the last few years, there are significant advances in the nonlinear self-trapping of weakly coupled condensates \cite{23,24}, long phase coherence time for two separated condensates \cite{25} etc. Stimulated by the remarkable experimental advances of BECs in double-well potential, we find strong parity-dependent interference fringes for ultracold bosonic gases released from a double-well potential, based on the phase diffusion in two-mode Bose-Hubbard model. For specific holding time of the double-well potential to create the phase diffusion, we find that the center in the density distribution is a dip for even particle number, while it is a peak for odd particle number. This provides a way to display directly the quantized characteristic of the particle number.

The paper is organized as follows. In Sec. II, we consider the phase diffusion of a condensate in a double-well potential, with the decoupled Bose-Hubbard Hamiltonian. The parity effect of the particle number is discussed based on the coherence property. In Sec. III, we predict strong parity effect of the particle number in the interference patterns of the density distribution. In Sec. IV, the influence of the parity effect due to asymmetry fluctuations of the double-well potential and particle number squeezing is studied. In the last section, we give a brief summary and discuss the application of Feshbach resonance to observe more clearly the parity effect.

II. PHASE DIFFUSION AND COHERENCE PROPERTY OF A CONDENSATE IN A DOUBLE-WELL POTENTIAL

For a condensate in a double-well potential with negligible tunneling, the system can be described by the following decoupled Bose-Hubbard Hamiltonian.

\begin{equation}
\hat{H} = \sum_{i=1}^{2} \varepsilon_i \hat{n}_i + \sum_{i=1}^{2} \frac{U}{2} \hat{n}_i (\hat{n}_i - 1). \tag{1}
\end{equation}
Here \( \tilde{n}_i = a_i^\dagger a_i \) is the particle number operator for the \( i \)th condensate. \( a_i \) and \( a_i^\dagger \) are bosonic annihilation and creation operators for the \( i \)th condensate. \( \varepsilon_1 - \varepsilon_2 \) represents the bias potential between two sites, while \( U \) denotes collisional interaction energy. The initial quantum state in this double-well potential is assumed as

\[
|\Psi (t_h = 0)\rangle = \frac{1}{\sqrt{2^N N!}} \left( a_1^\dagger + a_2^\dagger \right)^N |0\rangle. \tag{2}
\]

The evolution of the quantum state is then

\[
|\Psi (t_h)\rangle = \frac{1}{\sqrt{2^N N!}} e^{-\frac{i}{\hbar}H_{th}} \left( a_1^\dagger + a_2^\dagger \right)^N |0\rangle = \sum_{l=0}^{N} \frac{N!}{2^N N! (N-l)!} e^{-\frac{2iU th}{\hbar}} \left[ (l^{(l-1)}+(N-l)(N-l-1)) |l, N-l\rangle. \tag{3}\right.
\]

We see that there is a phase diffusion \[26, 27, 28, 29\] in the quantum state \(|\Psi (t_h)\rangle\) because of the exponential factor. In addition, from the exponential factor in \(|\Psi (t_h)\rangle\), the time evolution of \(|\Psi (t_h)\rangle\) displays a periodic behavior of \( T = 4\pi\hbar/U \). To give a clear discussion of the parity effect, in getting Eq. (2), we consider the case of \( \varepsilon_1 = \varepsilon_2 \).

The above evolution of \(|\Psi (t_h)\rangle\) may be realized by first preparing two coherently separated condensates in a double-well potential with large Josephson tunneling, so that the quantum state can be described well by Eq. (2). The central barrier is then increased non-adiabatically so that the tunneling between two condensates can be omitted, while the quantum state of the system still takes the form given by Eq. (2). This is similar to the experimental studies of the collapse and revival of the quantum coherence of the matter wave packet in an optical lattice \[30\]. This scheme was also applied in Ref. [4] to consider the macroscopic superpositions of phase states.

A direct way to display the coherence property of \(|\Psi (t_h)\rangle\) is to consider the time evolution of \( K(t_h) = \langle \Psi (t_h) | \left( a_1^\dagger a_2 + a_2^\dagger a_1 \right) |\Psi (t_h)\rangle /N \), which reflects the phase coherence and interference effect between two condensates. From the expression [3], we have

\[
K_{12} (t_h) = \frac{\langle \Psi (t_h) | \left( a_1^\dagger a_2 \right) |\Psi (t_h)\rangle}{N} = \sum_{l=1}^{N} \frac{(N-1)!}{2^N (l-1)! (N-l)!} e^{-\frac{2U th}{\hbar} (N-2l+1)}. \tag{4}\]

The above equation can be rewritten as

\[
K_{12} (t_h) = e^{-\frac{U t_h}{\hbar} (N-1)} \sum_{l=0}^{N-1} \frac{(N-1)!}{l! (N-l-1)!} \left( e^{\frac{2U th}{\hbar}} \right)^l. \tag{5}\]

Based on the binomial theorem, we have

\[
K_{12} (t_h) = \frac{e^{-\frac{U t_h}{\hbar} (N-1)}}{2^N} \left( 1 + e^{\frac{2U th}{\hbar}} \right)^{N-1}. \tag{6}\]

In this situation, we have the following simple expression

\[
K_{12} (t_h) = \frac{(\cos (U t_h/\hbar))^{N-1}}{2}. \tag{7}\]

From the above results, we have

\[
K(t_h) = (\cos (4\pi t_h))^{N-1}, \tag{8}\]

where \( t_h = t_h/T \). For large \( N \), by using the identity \( \lim_{N \to \infty} (1 - x/N)^N = e^{-x} \), the above expression can be approximated very well as

\[
K(t_h) \approx \sum_{n=0}^{N} (-1)^{n(N-1)} e^{-\frac{2U th}{\hbar} (4\pi t_h-n/4)^2}. \tag{9}\]

where \( n \) is an integer. In this situation, the width of the peaks or dips in \( K(t_h) \) is about \( 1/2\pi\sqrt{N} \). At \( t_h = 0 \), there is ideal phase coherence between two condensates. As shown in Fig. 1(a) for \( N = 101 \), with time increasing, the phase coherence disappears very rapidly for \( t_h > 1/2\pi\sqrt{N} \). With further time increasing, however, there is a revival of the phase coherence. This behavior is understood from the periodic behavior of the quantum state \(|\Psi (t_h)\rangle\).

A unique parity effect of the particle number is shown by the expression [5] of \( K(t_h) \). For specific holding time such that \( \cos (4\pi t_h) = -1 \), we see that \( K(t_h) = 1 \) for odd particle number, while \( K(t_h) = -1 \) for even particle number. This parity effect is shown further in Fig. 1(a) and Fig. 1(b) for \( N = 101 \) and \( N = 100 \), respectively.

III. PARITY EFFECT OF THE PARTICLE NUMBER IN THE DENSITY DISTRIBUTION

The above parity effect can be considered further by overlapping two condensates after switching off the
double-well potential. After a holding time of \( t_{\text{ho}} \), the quantum state is \( \ket{\Psi(t_{\text{ho}})} \). Then, we calculate the density distribution after a time of flight \( t_f \) by switching off the double-well potential. In our calculations, the initial wave functions of two condensates are assumed as \( \varphi_1(x) = e^{-\left(x-d/2\right)^2/2\sigma^2}/\pi^{1/4}\sigma^{1/2} \) and \( \varphi_2(x) = e^{-\left(x-d/2\right)^2/2\sigma^2}/\pi^{1/4}\sigma^{1/2} \), respectively. In this situation, the distance between two condensates is \( d \). In the following calculations, we will adopt the length unit \( d \), energy unit \( E_d = \hbar^2/2md^2 \) and time unit \( T_d = \hbar/E_d \). After a time of flight \( t_f = t_{f0}T_d \), the density distribution of the system is given by

\[
n(x, t_{\text{ho}}, t_f) = \frac{N}{2} \left[ \left| \phi_1(x, t_{f0}) \right|^2 + \left| \phi_2(x, t_{f0}) \right|^2 + 2 \times \text{Re} \left( (\cos(4\pi t_{\text{ho}}))^{N-1} \phi_1^* (x, t_{f0}) \phi_2 (x, t_{f0}) \right) \right]. \tag{10}
\]

For \( N = 100 \) and \( N = 101 \), in Figs. 2(a)-(f), we give the density distribution for fixed time of flight \( t_{f0} = 1 \) and different holding time \( t_{\text{ho}} \). In our calculations, \( \sigma = 0.1 \). The factor \( (\cos(4\pi t_{\text{ho}}))^{N-1} \) in the interference term of the above equation gives several unique results:

(i) There are collapse and revival of the interference fringes in the density distribution.

(ii) At \( t_{\text{ho}} = j/2 - 0.25 \) (\( j \) is a natural number), the density distribution depends strongly on the parity of the total particle number. As shown in Fig. 2(d), the center of the density distribution (solid line) is a dip for even particle number, while the center of the density distribution (dotted line) is a peak for odd particle number. This can be understood further by noticing that \( n = N \left| \phi_1 - \phi_2 \right|^2 /2 \) for even particle number, while \( n = N \left| \phi_1 + \phi_2 \right|^2 /2 \) for odd particle number. For both situations, the density distribution is symmetric, which means a clear parity effect about the total particle number \([31]\).

At \( t_{\text{ho}} = j/2 - 0.25 \), the factor \( \left( \cos(4\pi t_{\text{ho}}) \right)^{N-1} \) in the interference term can be rewritten as \( e^{i\pi(N-1)} \). In this situation, the phase diffusion of the left and right condensates is zero. However, there is a relative phase between two condensates. This relative phase is quantized about the total particle number. For even particle number, the relative phase is \( \pi \), while the relative phase is \( 2\pi \) for odd particle number.

Before the atomic cloud is imaged on a CCD camera, we do not know the parity of the particle number. At \( t_{\text{ho}} = j/2 - 0.25 \), the parity effect would lead to a random behavior in the density distribution. Averaging the density distribution in different experiments, there would be no interference patterns. However, in a single experiment, there will be clear interference patterns, and there is 50% probability to observe a peak (dip) in the center of the density distribution. This is significantly different from the holding time of \( t_{\text{ho}} = j/2 - 0.5 \), where one would always observe the same interference patterns in the density distribution.

These analyses at \( t_{\text{ho}} = j/2 - 0.25 \) can be considered further by the following density matrix

\[
\rho = P(N) \ket{N} \bra{N}. \tag{11}
\]

Here \( P(N) \) is the probability to have \( N \) atoms in the condensate. It can be assumed as a Gaussian distribution about the average particle number \( \overline{N} \). The above expression can be rewritten as

\[
\rho = \sum_{N=\text{even}} P(N) \ket{N} \bra{N} + \sum_{N=\text{odd}} P(N) \ket{N} \bra{N}. \tag{12}
\]

The ensemble average of the density distribution is then

\[
\overline{\rho}(x, t_{\text{ho}}, t_f) = \text{Tr} \left[ \rho \hat{\Theta} \right] = \frac{\overline{N}}{2} \left[ \left| \phi_1(x, t_{f0}) \right|^2 + \left| \phi_2(x, t_{f0}) \right|^2 \right]. \tag{13}
\]

It is clear that there is no interference term in the ensemble average of the density distribution. In a single-shot density distribution, however, there would be interference fringes with random relative phase \( \pi \) or \( 2\pi \).

**IV. INFLUENCE OF THE PARITY EFFECT DUE TO ASYMMETRY OF THE DOUBLE-WELL POTENTIAL AND PARTICLE NUMBER SQUEEZING**

As shown in preceding section, the density distribution depends strongly on the parity of the overall particle number \( N \). It is a natural problem whether slight asymmetry fluctuations of the double-well potential will destroy the parity effect. Generally speaking, the asymmetry fluctuations of the double-well potential will lead
to a random relative phase and fluctuations of the average particle number in two sites. To consider both effects, we study the following more general initial state

$$|\Psi (t_h = 0)\rangle = \frac{1}{\sqrt{2^N N!}} \left( \alpha \hat{a}_1^\dagger + \beta \hat{a}_2^\dagger \right)^N |0\rangle .$$  \hspace{1cm} (14)

The normalization condition requests that $|\alpha|^2 + |\beta|^2 = 1$. At a holding time $t_h$, the density distribution after a time of flight $t_f$ is given by

$$n(x, t_h, t_f) = \frac{N}{2} \left\{ |\alpha|^2 |\phi_1 (x, t_f)|^2 + |\beta|^2 |\phi_2 (x, t_f)|^2 \right\} + 2 \times \text{Re} \left[ \alpha^* \beta \frac{|\alpha|^2 e^{i4\pi t_h} + |\beta|^2 e^{-i4\pi t_h}}{2} \right]^{N-1} \phi_1^* (x, t_f) \phi_2 (x, t_f) \right\}. \hspace{1cm} (15)$$

At the holding time $t_h = j/2 - 0.25$, the density distribution is then

$$n(x, t_h, t_f) = \frac{N}{2} \left\{ |\alpha|^2 |\phi_1 (x, t_f)|^2 + |\beta|^2 |\phi_2 (x, t_f)|^2 \right\} + 2 \times (1)^{N-1} \times \text{Re} \left[ \alpha^* \beta \phi_1^* (x, t_f) \phi_2 (x, t_f) \right]. \hspace{1cm} (16)$$

We see that there is still strong parity effect about the total particle number when the asymmetry of the average particle numbers in two sites is considered. For $t_h = j/2 - 0.25$, $n = N|\alpha \phi_1 - \beta \phi_2|^2 / 2$ for even particle number, while $n = N|\alpha \phi_1 + \beta \phi_2|^2 / 2$ for odd particle number.

The factor $\alpha^* \beta$ includes a random relative phase if the fluctuations of the bias potential $\varepsilon_1 - \varepsilon_2$ are considered. The bias potential $\varepsilon_1 - \varepsilon_2$ between two sites will gives a factor $e^{-i(\varepsilon_1 - \varepsilon_2) t_h / \hbar}$. Experimentally, to observe the parity effect, this requests that the fluctuations of the bias potential should satisfy $|\Delta \varepsilon| t_h / \hbar << \pi$. This request can be satisfied in the present experimental technique [13, 23, 32]. These analyses show that slight asymmetry fluctuations of the double-well potential would not lead to serious problem in observing the parity effect.

In the non-adiabatical increasing of the central barrier so that the tunneling between two condensates can be omitted, it is possible that there would be a particle number squeezing. In this situation, we consider the following quantum state

$$|\Psi (t_h = 0)\rangle = \sum_{l=0}^{N} \frac{\gamma}{\sqrt{l! (N-l)!}} e^{-(l-N/2)^2 s^2 / N} \sqrt{l! (N-l)!} |l, N-l\rangle . \hspace{1cm} (17)$$

Here $s$ is the squeezing factor, and $\gamma$ is a normalization factor. For this quantum state, we have

$$K_{12} (t_h) = \frac{\gamma^2}{N} \sum_{l=1}^{N} \sqrt{l! (N-l+1)!} \times e^{-(l-N/2)^2 + (l-N/2-1)^2} s^2 / N \times e^{-\alpha t_h / 4(N-2l+1)} \times e^{-2l t_h / (N-2l+1)}. \hspace{1cm} (18)$$

At $t_h = j/2 - 0.25$, we have

$$K_{12} = \frac{\gamma^2}{N} \sum_{l=1}^{N} \sqrt{l! (N-l+1)!} \times e^{-(l-N/2)^2 + (l-N/2-1)^2} s^2 / N \times e^{-\alpha t_h / 4(N-2l+1)}. \hspace{1cm} (19)$$

At this holding time, we can also get

$$K_{21} = \frac{\gamma^2}{N} \sum_{l=0}^{N-1} \sqrt{(l+1)! (N-l)!} \times e^{-(l-N/2)^2 + (l-N/2+1)^2} s^2 / N \times e^{-\alpha t_h / 4(N-2l+1)}. \hspace{1cm} (20)$$

We see that at $t_h = j/2 - 0.25$, even particle number corresponds to a relative phase of $\pi$, while odd particle number corresponds to a relative phase of $2\pi$. This means that the squeezing factor will not destroy the parity effect in the interference fringes of the density distribution. Of course, there is a request that the quantum state should not be a Fock state in the extreme squeezing situation.

V. SUMMARY AND DISCUSSION

In summary, a strong parity effect of the total particle number in the density distribution of a condensate released from a double-well potential is predicted in this work. The asymmetry fluctuations of the double-well potential and particle number squeezing are discussed, and our studies show that they would not lead to serious problem in observing the parity effect. The parity effect is due to the quantized relative phase about the total particle number for specific holding time of the double-well potential. Another example of the strong parity effect about the total particle number is the exchange effect for a complex of $N$ spin 1/2 fermions [33]. The parity effect of the complex lies in that it is a boson for even $N$, while it is a fermion for odd $N$.

In this work, similar to the phase diffusion considered in Ref. [21], we ignore the phase diffusion in the increasing of the central barrier by assuming that the central barrier is increased fast enough, and we also ignore the phase diffusion in the time of flight. During the time of flight, the interaction energy per particle will decrease rapidly, and thus the phase diffusion may be ignored. In Ref. [4], the phase diffusion in the time of flight is also ignored. Nevertheless, overcoming these two phase diffusions would contribute to more clear observation of the parity effect. One way to overcome these two phase diffusions would be the application of Feshbach resonance which has been used in atom interferometry [34, 35]. The scheme with the application of Feshbach resonance would be: (i) One first prepares two coherently separated condensates in a double-well potential. (ii) After turning the s-wave scattering length almost to zero via
a magnetic-field Feshbach resonance, the central barrier is increased adiabatically so that the tunneling between two sites can be omitted. (iii) Holding the double-well potential, one adiabatically increases and then decreases the s-wave scattering length almost to zero so that the condition \( \int_0^t U(t) dt = \pi \) is satisfied. (iv) Switching off the double-well potential, the ideal gas is then imaged after a time of flight to display the parity effect.

**Acknowledgments**

We acknowledge useful discussions with Prof. Y. Wu, Prof. B. Wu, Prof. B. L. Lv and Prof. L. You. This work was supported by NSFC under Grant Nos. 10875165, 10804123, 10634060, and NKBRSF of China under Grant No. 2006CB921406.

[1] N. Hatakenaka, Phys. Rev. Lett. 81, 3753 (1998).
[2] Z. D. Chen, J. Q. Liang, S. Q. Shen, and W. F. Xie, Phys. Rev. A 69, 023611 (2004).
[3] R. Lü, M. Zhang, J. L. Zhu, and L. You, Phys. Rev. A 78, 011605(R) (2008).
[4] F. Piazza, L. Pezzè, and A. Smerzi, Phys. Rev. A 78, 051601(R) (2008).
[5] V. S. Shchesnovich, Preprint arXiv: 0905.1708v2 (2009).
[6] F. D. M. Haldane, Phys. Rev. Lett. 90, 1029 (1988).
[7] D. Loss, D. P. DiVincenzo, and G. Grinstein, Phys. Rev. Lett. 69, 3232 (1992).
[8] Y. Shin, M. Saba, T. A. Pasquini, W. Ketterle, D. E. Pritchard, and M. Prentiss, Phys. Rev. Lett. 98, 030407 (2007).
[9] E. M. Wright, D. F. Walls, and J. C. Garrison, Phys. Rev. Lett. 77, 2158 (1996).
[10] J. von Delft and C. L. Henley, Phys. Rev. Lett. 69, 3236 (1992).
[11] Y. Shin, M. Saba, T. A. Pasquini, W. Ketterle, D. E. Pritchard, and A. E. Leanhardt, Phys. Rev. Lett. 92, 050405 (2004).
[12] Y. Shin, G. B. Jo, M. Saba, T. A. Pasquini, W. Ketterle, and D. E. Pritchard, Phys. Rev. Lett. 95, 170402 (2005).
[13] T. Schumm, S. Hofferberth, L. M. Andersson, S. Wildermuth, S. Groth, J. Bar-Joseph, J. Schmiedmayer, and P. Kruger, Nature Physics 1, 57 (2005).
[14] M. Greiner, O. Mandel, T. W. Hansch, and I. Bloch, Nature 419, 51 (2002).
[15] M. Greiner, O. Mandel, T. W. Hansch, and I. Bloch, Nature 419, 51 (2002).
[16] Y. Shin, G. B. Jo, M. Saba, T. A. Pasquini, W. Ketterle, and D. E. Pritchard, Phys. Rev. Lett. 95, 170402 (2005).
[17] Y. Shin, M. Saba, T. A. Pasquini, W. Ketterle, and D. E. Pritchard, Phys. Rev. Lett. 95, 170402 (2005).
[18] M. Greiner, O. Mandel, T. W. Hansch, and I. Bloch, Nature 419, 51 (2002).
[19] Y. Shin, M. Saba, T. A. Pasquini, W. Ketterle, and D. E. Pritchard, Phys. Rev. Lett. 95, 170402 (2005).
[20] Y. Shin, M. Saba, T. A. Pasquini, W. Ketterle, and D. E. Pritchard, Phys. Rev. Lett. 95, 170402 (2005).
[21] Y. Shin, M. Saba, T. A. Pasquini, W. Ketterle, and D. E. Pritchard, Phys. Rev. Lett. 95, 170402 (2005).
[22] Y. Shin, M. Saba, T. A. Pasquini, W. Ketterle, and D. E. Pritchard, Phys. Rev. Lett. 95, 170402 (2005).
[23] Y. Shin, M. Saba, T. A. Pasquini, W. Ketterle, and D. E. Pritchard, Phys. Rev. Lett. 95, 170402 (2005).
[24] Y. Shin, M. Saba, T. A. Pasquini, W. Ketterle, and D. E. Pritchard, Phys. Rev. Lett. 95, 170402 (2005).
[25] Y. Shin, M. Saba, T. A. Pasquini, W. Ketterle, and D. E. Pritchard, Phys. Rev. Lett. 95, 170402 (2005).
[26] Y. Shin, M. Saba, T. A. Pasquini, W. Ketterle, and D. E. Pritchard, Phys. Rev. Lett. 95, 170402 (2005).
[27] Y. Shin, M. Saba, T. A. Pasquini, W. Ketterle, and D. E. Pritchard, Phys. Rev. Lett. 95, 170402 (2005).
[28] Y. Shin, M. Saba, T. A. Pasquini, W. Ketterle, and D. E. Pritchard, Phys. Rev. Lett. 95, 170402 (2005).
[29] Y. Shin, M. Saba, T. A. Pasquini, W. Ketterle, and D. E. Pritchard, Phys. Rev. Lett. 95, 170402 (2005).
[30] Y. Shin, M. Saba, T. A. Pasquini, W. Ketterle, and D. E. Pritchard, Phys. Rev. Lett. 95, 170402 (2005).
[31] Y. Shin, M. Saba, T. A. Pasquini, W. Ketterle, and D. E. Pritchard, Phys. Rev. Lett. 95, 170402 (2005).
[32] Y. Shin, M. Saba, T. A. Pasquini, W. Ketterle, and D. E. Pritchard, Phys. Rev. Lett. 95, 170402 (2005).
[33] Y. Shin, M. Saba, T. A. Pasquini, W. Ketterle, and D. E. Pritchard, Phys. Rev. Lett. 95, 170402 (2005).
[34] Y. Shin, M. Saba, T. A. Pasquini, W. Ketterle, and D. E. Pritchard, Phys. Rev. Lett. 95, 170402 (2005).
[35] Y. Shin, M. Saba, T. A. Pasquini, W. Ketterle, and D. E. Pritchard, Phys. Rev. Lett. 95, 170402 (2005).