Steering Bose-Einstein condensates despite time symmetry

Dario Poletti, Giuliano Benenti, Giulio Casati, Peter Hänggi, and Baowen Li

1Department of Physics and Centre for Computational Science and Engineering, National University of Singapore, Singapore 117542, Republic of Singapore
2CNISM, CNR-INFM, and Center for Nonlinear and Complex Systems, Università degli Studi dell’Insubria, Via Valleggio 11, 22100 Como, Italy
3Istituto Nazionale di Fisica Nucleare, Sezione di Milano, Via Celoria 16, 20133 Milano, Italy
4Institut für Physik, Universität Augsburg, Universitätsstr. 1, D-86135 Augsburg, Germany
5NUS Graduate School for Integrative Sciences and Engineering, 117597, Republic of Singapore

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A Bose-Einstein condensate in an oscillating spatially asymmetric potential is shown to exhibit a directed current for unbiased initial conditions despite time symmetry. This phenomenon occurs only if the interaction between atoms, treated in mean-field approximation, exceeds a critical value. Our findings can be described with a three-mode model (TMM). These TMM results corroborate well with a many-body study over a time scale which increases with increasing atom number. The duration of this time scale probes the validity of the used mean-field approximation.

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The realization of Bose-Einstein condensates (BECs) of dilute atomic gases has opened new possibilities for the investigation of interesting physical phenomena induced by atom-atom interactions. For example, it was found that for a BEC in a double-well potential the tunneling can be suppressed due to interaction-induced self-trapping, thus maintaining population imbalance between the two wells. A similar phenomenon has been predicted also for a periodically driven BEC in a double well and shows the importance of interaction in the coherent control of quantum tunneling between wells.

In recent years, ample interest has arisen in the field of directed transport in classical and quantum systems in absence of a dc-bias, see and Refs. therein. In particular it has been shown that by breaking space- and time-inversion symmetries it is possible to achieve directed (ratchet) transport. The study of the role of interactions on the quantum ratchet transport, however, truly remains terra incognita, with only a very few preliminary studies available. It has been shown that atom-atom interactions in a BEC can cause a finite current while the non-interacting system would not exhibit directed transport; however, all these systems possess a Hamiltonian that is both space- and time-asymmetric.

With this study, instead, we consider a non-dissipative, i.e. Hamiltonian system in which the time-reversal symmetry is not broken and the system is driven smoothly. We show that, starting out from initial conditions symmetric both in momentum and space, interactions, treated in the limits of a mean-field approximation, induce directed current in a BEC. In a full many-body analysis of the system, symmetry considerations imply that this directed current is asymptotically decaying. Despite this fact we show that a finite current persists over a time scale which increases with increasing atom number in the condensate.

To start out, we consider a BEC composed of N atoms confined in a toroidal trap of radius R and cross section πr², subjected to the condition r ≪ R, so that the motion is essentially one-dimensional. The condensate is driven by a time-periodic potential, reading

\[ V_{\text{Ext}}(x, t) = [V_1 \cos(x) + V_2 \cos(2x + \phi)] \cos[\omega(t - t_0)], \]

(1)

where \( V_1 \) and \( V_2 \) denote the potential depths of the two lattice harmonics, \( \phi \) their relative phase, and \( \omega \) the angular frequency of the oscillations of the perturbation. For a trap of radius \( R = 2.5 \mu m \), our used unit of time is \( mR^2/\hbar = 6.93 \times 10^{-4} s \). This potential can be readily reproduced in experiments. The space symmetry \( \{x \rightarrow -x + \gamma, \text{with } \gamma \text{ constant} \} \) is broken when \( \phi \neq 0, \pi \). The driving is symmetric under \( t \rightarrow -t + 2t_0 \) inversion and in particular, for \( t_0 = 0 \), under \( t \rightarrow -t \) time symmetry.

At zero temperature and as long as the number \( N_0 \) of condensate particles is much larger than the number \( \delta N \) of non-condensate ones, the evolution of a BEC is well described by the Gross-Pitaevskii (GP) equation. Thus, for our system the evolution of the condensate wave-function \( \psi(x, t) \), normalized to 1, is described by (taking \( \hbar = m = 1 \))

\[ i \frac{\partial \psi(x, t)}{\partial t} = \left[ -\frac{1}{2} \frac{\partial^2}{\partial x^2} + g|\psi(x, t)|^2 + V_{\text{Ext}}(x, t) \right] \psi(x, t), \]

(2)

where \( x \in [0, 2\pi] \) is the azimuthal angle of the torus, \( g = 8NaR/\mu \) the scaled strength of the nonlinear interaction (we consider the repulsive case, i.e., \( g > 0 \)) and \( a \) is the s-wave scattering length for elastic atom-atom collisions.

As initial condition we use a wave-function which is symmetric both in space, \( x \), and in momentum, \( p \), i.e. it
is of the form \( \psi(x,0) = \sum_n a_n e^{i n x} \cos(n \pi x) \), where \( n \in \mathbb{N} \) and \( a_n, \alpha_n \in \mathbb{R} \). For simplicity, and without loss of generality we choose the initial condition to be:

\[
\psi(x,0) = \frac{1}{\sqrt{2}} \left[ \frac{1}{\sqrt{2\pi}} + e^{i \alpha} \cos(x) \right],
\]

where \( \alpha \) is the relative phase between the homogeneous part and the last term which corresponds to a cosinusoidal modulation in the initial density of particles. While in the non-interacting \((g = 0)\) case the asymptotic current is zero for any initial condition that is symmetric in \( x \) and \( p \), we show that this is not the case if \( g \neq 0 \).

As depicted with Fig. 1 (top panel), our model may exhibit a non-zero asymptotic current, \( \langle p \rangle_{\text{asym}} = \lim_{t \to \infty} \frac{1}{\pi} \int_0^{\pi} \langle p \rangle(t') dt' \), when \( g \neq 0 \) and \( \alpha \neq 0 \) with \( \langle p \rangle(t') = -i \int_0^{2\pi} \psi^*(x,t') \frac{\partial}{\partial x} \psi(x,t') dx \). The asymptotic time-averaged momentum \( \langle p \rangle_{\text{asym}} \) vs. \( g \), Fig. 1 (bottom panel), clearly shows that (i) there is a finite threshold value \( g^* \) such that \( \langle p \rangle_{\text{asym}} \neq 0 \) when \( g > g^* \) and (ii) there exists an optimal value \( g_{\text{opt}} \) that maximizes the current.

The asymptotic time-averaged current is shown as a function of the initial phase \( \omega t_0 \) of the driving field in Fig. 2. We can deduce that there is a strong asymmetry around \( \omega t_0 = \pi \), in particular for the coupling strength \( g = 0.075 \). As a result, the directed current survives even after averaging over \( t_0 \) (cf. the inset of Fig. 2).

![FIG. 1: (color online) Upper panel: Momentum expectation \( \langle p \rangle \) vs. time \( t \) for \( \alpha = 0 \) (dark black curve) and \( \alpha = \pi/2 \) (light red curve). Used parameter values are: \( g = 0.2, V_1 = V_2 = 2, \phi = \pi/2, \omega = 10, t_0 = 0 \). Lower panel: Asymptotic time-average momentum \( \langle p \rangle_{\text{asym}} \) versus the interaction strength \( g \) for \( \alpha = \pi/2 \). Data are obtained from the GP equation (2) (black squares) or from the TMM-Ansatz (5) (red circles). In the first case \( g^* \approx 0.065 \) and \( g_{\text{opt}} \approx 0.15 \), in the latter \( g^* \approx 0.070 \) and \( g_{\text{opt}} \approx 0.15 \).](image)

![FIG. 2: (Color online) Asymptotic time-averaged momentum \( \langle p \rangle_{\text{asym}} \) versus \( \omega t_0 \) for \( g = 0.05 \) (filled black squares), \( g = 0.075 \) (empty red circles), \( g = 0.1 \) (filled blue triangles), and \( g = 0.2 \) (pink asterisks). Inset: asymptotic current averaged over \( t_0 \), \( \langle p \rangle_{\text{asym}} \), as a function of the interaction strength \( g \). Parameter values: \( V_1 = V_2 = 2, \phi = \pi/2, \omega = 10, \alpha = \pi/2 \).](image)
the maximum symmetry violation values $\phi = \pi/2, 3\pi/2$.

From our numerical studies we notice that the only states that are significantly excited during the dynamical evolution of the system assume the momenta 0, +1 or −1. We hence attempt to reproduce our present results, at least on a qualitative level, with the use of a mean-field three-mode model (TMM). We use the Ansatz

$$\psi(x, t) = \frac{1}{\sqrt{2\pi}} \left[ A(t) + P(t)e^{ix} + M(t)e^{-ix} \right],$$

where $A(t), P(t)$ and $M(t)$ are time-dependent, complex coefficient such that $|A|^2 + |P|^2 + |M|^2 = 1$. Using (4) we find the effective mean-field Hamiltonian of the TMM:

$$H_{\text{eff}} = \int_0^{2\pi} dx \psi^*(x, t) \left[ -\frac{1}{2} \frac{\partial^2}{\partial x^2} + \frac{g}{2N} \frac{v}{\pi^2} \psi^2 + V_{\text{Ext}}(x, t) \right] \psi,$$

$$+ \frac{K}{2} (P + M) \cos(\omega t),$$

$$+ \frac{g}{2} A^2 M \cos(\omega t).$$

(6)

This TMM qualitatively reproduces the behavior of the asymptotic time-averaged current, see Fig. 3 (bottom panel). In Fig. 4 we depict the evolution of the populations $|A|^2$, $|P|^2$, and $|M|^2$ of the three modes. For $\alpha = 0$ ($\alpha = \pi/2$) the time-averaged current is zero (non-zero).

The obtained results may appear surprising if one observes that the full second-quantized quantum many-body operator is linear, obeying time-reversal symmetry, and symmetry arguments necessarily imply a vanishing asymptotic current [14]. Actually there occurs no contradiction because the two limits $t \to \infty$ and $N \to \infty$ do not commute: a directed current is obtained only in the mean-field limit in which we let first $N \to \infty$ and then $t \to \infty$. The issue for a physical system in which both the number of particles and the time of the experimental run are finite constitutes the vindication of the regime of validity of the used nonlinear GP equation. For this purpose, we next study the second quantized many-body problem within the three-mode approximation, being also accessible to numerical investigation.

The second-quantized quantum many-body Hamiltonian reads:

$$\hat{H}_{\text{eff}} = \int \frac{dx}{2\pi} \psi^\dagger \left[ -\frac{1}{2} \frac{\partial^2}{\partial x^2} + \frac{g}{2N} \frac{v}{\pi^2} \psi^\dagger \psi + V_{\text{Ext}}(x, t) \right] \psi,$$

where, within the three-mode approximation,

$$\psi^\dagger = \frac{1}{\sqrt{2\pi}} \left( \hat{a}^\dagger + \hat{p}^\dagger e^{ix} + \hat{m}^\dagger e^{-ix} \right),$$

(8)

(9)

$\hat{a}^\dagger, \hat{p}^\dagger$, and $\hat{m}^\dagger$ denote the Boson creation operators for the states with momenta 0, +1, and −1. These operators satisfy the commutation relations $[\hat{a}, \hat{a}^\dagger] = 1,$ $[\hat{p}, \hat{p}^\dagger] = 1,$ $[\hat{m}, \hat{m}^\dagger] = 1$ and operators corresponding to different modes commute. The particle conservation reads $(\hat{a}^\dagger \hat{a} + \hat{p}^\dagger \hat{p} + \hat{m}^\dagger \hat{m}) = N$, where $N$ is the total number of particles. As initial condition, we consider $N$
atoms in the state
\[
\frac{1}{\sqrt{N!}} \left( \frac{\hat{a}^\dagger}{\sqrt{2}} + \frac{e^{i\alpha} (\hat{p}^\dagger + \hat{m}^\dagger)}{2} \right)^N |0\rangle.
\]  

(10)

We studied numerically the evolution of the \( N \)-body system governed by the second-quantized three-mode Hamiltonian [5]. The normalized population imbalance between levels with momenta \(+1\) and \(-1\) is depicted in Fig. 5. The TMM holds up to increasingly longer times \( t^* \) as \( N \) increases. We estimate \( t^*(N) \) as the time instant at which the \( N \)-body population imbalance deviates by more than 1% from the mean-field treatment value: Our numerical data are well fitted by a logarithmic dependence, \( t^* \propto \ln N \), cf. inset of Fig. 5.

To check the stability of the condensate, we have diagonalized the single-particle reduced density matrix, whose matrix elements read \( \langle \hat{x}^m \hat{y}^n \rangle \), with \( \hat{x}, \hat{y} \) being \( \hat{a}, \hat{p} \) or \( \hat{m} \). According to [10], we have a BEC if this matrix possesses one eigenvalue of order \( N \). Defining \( \tilde{t} \) as the time for which the condensed state is highly populated, we have found numerically that \( \tilde{t} \propto N^{0.45} \). This means that for large \( N \) and times \( t \) obeying \( \tilde{t} > t > t^* \) the mean-field approach no longer provides an accurate description of the many-body system despite the presence of a condensate.

In conclusion, we have shown that in BEC the interaction plays a prominent role in inducing long lasting currents, when otherwise impossible. We note that BECs in toroidal traps have recently been realized in different experiments [21, 21, 22]... The initial condition [3] can be implemented by properly exciting the ground-state of a condensate in a toroidal trap, and the phase \( \alpha \) can be tuned by letting the excited condensate evolve over a given time span. The interaction strength can be changed by tuning an external magnetic field [24]. Extrapolating the dependence given in Fig. 5 of \( t^* \) on the number \( N \) of condensate particles, we obtain \( t^* \approx 800 \) (in an in situ experiment this corresponds a time scale 0.1 s) for \( N \approx 5 \times 10^5 \) particles [22]. Changing the number of condensed particles [21] opens the possibility to explore in situ the range of validity of the GP description via the measurement of our predicted, intriguing directed transport.

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\[ \begin{align*}
\text{FIG. 5: (Color online) Population imbalance divided by } N \text{ between levels with } p = +1 \text{ and } p = -1 \text{ vs. } t. \text{ From bottom to top: } N = 10, 80, \text{ and mean-field TMM, for } g = 0.2, V_1 = V_2 = 2, \phi = \pi/2, \omega = 10, t_0 = 0, \alpha = \pi/2. \text{ Inset: } t^* \text{ versus } N (\text{squarees}) \text{ and numerical fit } t^* = A + B \ln N, \text{ with } A \approx 73 \text{ and } B \approx 54.
\end{align*} \]

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