Two-dimensional Bose-Einstein condensate under extreme rotation

Sinha, S.; Shlyapnikov, G.V.

Published in:
Physical Review Letters

Citation for published version (APA):
Two-Dimensional Bose-Einstein Condensate under Extreme Rotation

S. Sinha\textsuperscript{1} and G. V. Shlyapnikov\textsuperscript{2,3}

\textsuperscript{1}Max-Planck-Institut für Physik Komplexer Systeme, 38 Nöthnitzer Straße, 01187 Dresden, Germany
\textsuperscript{2}Laboratoire Physique Théorique et Modèles Statistiques, Bâtiment 100, Université Paris-Sud, 91405 Orsay, France
\textsuperscript{3}Van der Waals-Zeeman Institute, University of Amsterdam, Valckenierstraat 65/67, 1018 XE Amsterdam, The Netherlands

(Received 23 December 2004; published 22 April 2005)

We show that a Bose-condensed gas, under extreme rotation in a 2D anisotropic trap, forms a novel elongated quantum fluid which has a roton-maxon excitation spectrum. For a sufficiently large interaction strength, the roton energy reaches zero and the system undergoes a second order quantum transition to the state with a periodic structure—rows of vortices. The number of rows increases with the interaction, and the vortices eventually form a triangular Abrikosov lattice.

\begin{equation}
H = \int d^2r \Psi^\dagger \left[ \frac{-i\hbar\nabla + m\vec{r} \times \vec{\Omega}}{2m} + V_{\text{ext}}(\vec{r}) + \frac{g}{2} \Psi^\dagger \Psi \right] \Psi,
\end{equation}

where $\Psi^\dagger$ and $\Psi$ are bosonic field operators, $m$ is the particle mass, $g$ is the coupling constant for the mean-field interaction, and the effective trapping potential is $V_{\text{eff}}(\vec{r}) = (m/2)[(\omega_x^2 - \Omega^2)x^2 + (\omega_y^2 - \Omega^2)y^2]$. The Hamiltonian (1) is analogous to that of charged particles in the magnetic field, and in this respect the quantity $m\vec{r} \times \vec{\Omega}$ is the gauge field. We consider the limit of extreme rotation, where $\Omega = \omega_x = \omega_y \sqrt{1 - \epsilon}$ and the gas becomes free along the $x$ direction. Then, assuming that in this direction the atoms are confined in a large rectangular box of size $L$, the gas becomes a long cigar. After the gauge transformation $\Psi = \Psi e^{im\Omega(x+y)/\hbar}$, the Hamiltonian in the Landau gauge can be written as

\begin{equation}
H = \int d^2r \Psi^\dagger \left[ \frac{(p_x + 2m\Omega y)^2 + p_y^2}{2m} + \frac{m\omega_+^2 y^2}{2} \right. \\
+ \left. \frac{g}{2} \Psi^\dagger \Psi \right] \Psi,
\end{equation}

where $\omega_+ = \omega_y \sqrt{2 \epsilon} \ll \Omega$, assuming a small ellipticity $\epsilon$.

Omitting the interaction term, the single particle eigenstates are the Landau levels [11] separated from each other by an energy gap $\sim 2\hbar \Omega$. In the dilute limit, where the mean-field interaction $g n_{2D} \ll \hbar \Omega$ ($n_{2D}$ is the 2D density), we may restrict our discussion within the lowest Landau level. Then the field operator can be written in the form $\Psi = \sum_k \phi_k a_k$, where $a_k$ is the creation operator of a particle with momentum $k$ along the $x$ direction and $\phi_k$ is the corresponding eigenfunction:

\begin{equation}
\phi_k(x, y) = \left( \frac{\pi \hbar L^2}{2} \right)^{1/4} \exp \left[ -\frac{1}{2} \left( \frac{y}{l_0} + \frac{\Omega k_0}{\tilde{\Omega}} \right)^2 \right],
\end{equation}

where $\tilde{\Omega} = \sqrt{\Omega^2 + \omega_+^2}/4$, and $l_0 = (\hbar/2m\tilde{\Omega})^{1/2}$. Then,
after the spatial integration of Eq. (2), we obtain an effective one-dimensional (1D) Hamiltonian:

\[
H = \sum_k (\hbar^2 k^2 / 2m^*) a_k^\dagger a_k + (g^* / 2L) \sum_{k'q} a_{k+q}^\dagger a_{k-q}^\dagger a_k a_k \exp(-l_0^2 (k - k' + q)^2) + q^2 / 2),
\]

(4)

where \( g^* = g / \sqrt{2 \pi l_0} \) is an effective 1D coupling constant. The Hamiltonian (4) describes particles with a large effective mass \( m^* = m(2\Omega / \omega_n)^2 \gg m \). The fact that \( m^* \) is not infinite and the kinetic energy term is still present originates from the asymmetry of the trapping potential. This asymmetry leads to a small difference between the frequencies \( \Omega \) and \( \Omega \). However, once the finite kinetic energy term is extracted one may put \( \Omega = \Omega \), and we have done this in the second term on the right-hand side of Eq. (4).

The momentum dependence of the interaction term in the Hamiltonian (4) originates from the presence of the gauge field. The wave functions of particles which have opposite momenta in the \( x \) direction are shifted in opposite directions along the \( y \) axis, which decreases their overlap and reduces the interaction amplitude.

The behavior of the system is governed by the particles with momenta \( k \ll l_0^{-1} \), for which the extension of the wave function in the \( y \) direction is \( l_0 \). If the 1D density \( n = N / L \) (\( N \) is the total number of particles) satisfies the condition \( nl_0 \ll 1 \), then we are dealing with a 1D Bose gas. In this case, characteristic momenta that are important are of the order of \( n \) or smaller. They satisfy the condition \( kl_0 \ll 1 \), and the exponential term in Eq. (4) is equal to unity. Then the Hamiltonian (4) corresponds to the well-described Lieb-Liniger model for the 1D Bose gas [12].

We, therefore, focus on the other extreme, where

\[
nl_0 \gg 1. 
\]

(5)

Then the system can be viewed as a 2D gas in a narrow channel. The 2D density is \( \sim n / l_0 \) and the interaction energy per particle is \( I \sim ng / l_0 \). The characteristic kinetic energy of a particle at the mean distance from other particles is \( K \sim h^2 n / m^* l_0 \), and for \( K \gg I \) the wave function of the particle at such interparticle distances is not influenced by the interactions. The gas is then in the weakly interacting mean-field regime. The criterion of weak interactions takes the form (see, for example, [13])

\[
m^* g / h^2 \ll 1. 
\]

(6)

Thus, under the conditions (5) and (6) we have a weakly interacting 2D Bose gas in a narrow channel.

As well as in the 1D Bose gas, the ground state can be a quasicondensate in which the density fluctuations are suppressed, but the phase fluctuates in the \( x \) direction on a distance scale smaller than the size \( L \) [13]. However, locally the quasicondensate is indistinguishable from a true BEC. The analysis in the density-phase representation [13] gives the same excitation spectrum as the one obtained in the Bogoliubov approach assuming that most particles are in the condensate. Employing this approach, we first reduce the Hamiltonian (4) to a bilinear form:

\[
H_b = \sum_k (\hbar^2 k^2 / 2m^* + 2n g^* \exp(-k^2 l_0^2 / 2)) a_k^\dagger a_k + (ng^* / 2) \sum_k \exp(-k^2 l_0^2 / 2)(a_k^\dagger a_k + a_k a_k). 
\]

(7)

Diagonalizing \( H_b \) we obtain the excitation spectrum:

\[
e^2(k) = [\hbar^2 k^2 / 2m^* + g^* n (2 \exp(-k^2 l_0^2 / 2) - 1)]^2 - g^* n^2 \exp(-2k^2 l_0^2).
\]

(8)

The key feature of the excitation energy (8) is the momentum dependence of the interaction terms proportional to \( g^* \). The structure of the spectrum depends on the ratio of the mean-field interaction to the kinetic energy at momentum \( k = 1 / l_0 \). This ratio can take both small and large values and is given by

\[
\beta = n g^* / (2 \pi^2 h^2) = \sqrt{2} (m^* g / \hbar^2) n l_0.
\]

(9)

In units of \( mh^2 / m^* \), the excitation energy is a universal function of \( \beta \) and \( kl_0 \). For small \( \beta \), the interaction terms proportional to \( g^* \) in Eq. (8) are important only at \( k \ll 1 / l_0 \), where they become momentum independent. Then Eq. (8) gives the ordinary Bogoliubov spectrum, with a small sound velocity \( c_s = \sqrt{ng^*/m^*} \).

For \( \beta \gg 1 \) the situation drastically changes. The interaction is already important for momenta \( k \geq 1 / l_0 \), where the interaction dependent terms in Eq. (8) decrease with increasing \( k \). For this reason the spectrum develops the roton-maxon structure for \( \beta > 2.6 \) (see Fig. 1). This structure is well known in the physics of superfluid \(^4\text{He}\) and is predicted for trapped dipolar condensates [14] and studied for lattice bosons with long-range interactions [15]. The roton minimum is located at \( k = k_c = 1.6 / l_0 \), and the corresponding excitation energy decreases with increasing \( \beta \). For a critical value \( \beta = 4.9 \), the roton minimum reaches zero at \( k = k_c = 1.6 / l_0 \), and a further increase in \( \beta \) makes the

FIG. 1. Excitation energy (in units of \( mh^2 / m^* \)) versus \( kl_0 \).
The absence of current in the form of this type of wave function reads
\[ \psi = \sqrt{N} \left[ C_0 \phi_0 + \sum_{j=1}^{j} (C_k \phi_{k} + C_{-k} \phi_{-k}) \right]. \tag{10} \]

The energy of the 3-component state is
\[ E/N = g^* n/2 - B^2/4A, \] when the single to 3-component wave function can be treated as a second order quantum transition. The energy and chemical potential \( \mu \) change continuously, whereas the quantity \( \partial\mu/\partial n \) undergoes a jump at the transition point. For the 3-component state it is smaller by an amount \( 0.5g^* \).

Near the transition point quantum fluctuations increase due to the vanishing excitation energy at the roton minimum. This energy can be expressed as
\[ \varepsilon(k) = \left( \frac{\hbar^2}{2m^* l_0^2} \right) \frac{2}{\beta_e} (\beta_e - \beta) + \gamma l_0^4 (k^2 - k_0^2)^2 \right]^{1/2}, \]
where \( \gamma = 0.12 \). For the mean square fluctuations of the density we then obtain
\[ \left( \frac{\delta n}{n} \right)^2 = \int \frac{dk}{2n\pi} \left[ \frac{\hbar^2 k^2}{2m^* \varepsilon(k)} - 1 \right] = \frac{1}{n l_0} \ln \left( \frac{1}{\beta_e - \beta} \right). \]

These fluctuations become large for \( \delta \beta = (\beta_e - \beta) \leq \exp(-n l_0) \). Thus, the transition from a single to 3-component wave function occurs in the interval \( \delta \beta \), which is exponentially narrow due to the inequality (5). The behavior of the system in this region requires a special investigation and is beyond the scope of this Letter.

It is important that already for \( \beta = 5.4 \) the ground state wave function changes from 3- to 2-component, and it again becomes 3-component at \( \beta = 20 \). With increasing \( \beta \) (either increasing \( g^* n \) or decreasing \( \omega_\perp \)), more momentum states are macroscopically occupied. This is because an increase of \( \beta \) is equivalent to increasing the effective mass \( m^* \), which makes momentum states in the lowest Landau level (LLL) more degenerate.

Our results for the number of momentum components in the ground state wave function are displayed in Fig. 2. These states describe one or several vortex rows along the \( x \) axis. For example, a two component state represents one vortex row [see Fig. 3(I)]. Eventually, a triangular Abrikosov vortex lattice [17] is formed when increasing \( \beta \) to very large values. It should be mentioned here that the structure of vortex rows has been obtained in the studies of type-II superconductors [18] and in the studies of condensates in rotating anisotropic traps [19].

FIG. 2. The number of components \( (N_r) \) in the ground state macroscopic wave function for various values of \( \beta \).
In the presence of many rows of vortices, we may rely on an average vortex description in the LLL. Using the relation \( \langle \phi_k | y^2 | \phi_k \rangle = k^2 l_0^2 + 1/2 \), we write the energy functional of the system in the form \( E[\rho] = \int d^2 r \{ m \omega^2 y^2 \rho + g \rho^2 \} / 2 \). Minimizing the energy, we obtain the coarse grained density \( \rho(y) = \frac{1}{g} [\mu - \frac{1}{2} m \omega^2 y^2] \), which smooths out density modulations introduced by vortex rows. The size of the vortex core is always \( \sim l_0 \), and the number of the vortex rows is \( \sim \beta^{1/3} \). The energy per particle is given by

\[
E/N = 3 \hbar \Omega (\omega_0 / 2 \Omega)^2 (3\sqrt{2} \pi \beta / 4)^{2/3} / 5. \tag{15}
\]

The result of Eq. (15) deviates by less than 10% from the energy obtained by the full numerical minimization.

The vortex lattice is expected to melt when the number of vortices \( N_v \) approaches the number of particles \( N \) [3,20]. In our case the condition \( N_v \sim N \) transforms into \( m^* g / \hbar^2 \sim n^2 l_0^2 \). However, our approach of the weakly interacting mean-field regime in a narrow channel breaks down before the melting transition. This approach requires inequalities (5) and (6) and, hence, is surely expected to be invalid for \( m^* g / \hbar^2 \gtrsim n^2 l_0^2 \). Therefore, the description of the melting transition requires a more elaborate treatment of quantum fluctuations.

In conclusion, we have shown that a 2D BEC at extreme rotation frequency in an elliptic trap forms a novel quantum fluid in a narrow channel. The behavior of this fluid is determined by the parameter \( \beta \), Eq. (9) which increases with the interaction strength and effective mass \( m^* \). For \( \beta < \beta_c = 4.9 \), the excitation spectrum of the fluid has a small sound velocity and exhibits a roton-maxon character. At a critical interaction strength \( \beta_c \), the roton energy reaches zero and the uniform (in the long direction) ground state becomes unstable. The system then undergoes a second order quantum transition to the state with a periodic structure which can be viewed as rows of vortices. For \( \beta > \beta_c \), with increasing the interaction parameter \( \beta \), more vortex rows are nucleated. Finally, the vortices form the Abrikosov lattice which can melt due to quantum fluctuations.

We acknowledge discussions with J. Dalibard, A. L. Fetter, Tin-Lun Ho, J. Palacios, and J.T.M. Walraven. This work was supported by the Ministère de la Recherche (Grant ACI Nanoscience 201), by the Centre National de la Recherche Scientifique (CNRS), by the Nederlandse Stichting voor Fundamenteel Onderzoek der Materie (FOM), and in part by the National Science Foundation (Grant No. PHY99-07949), LPTMS is a mixed research unit of CNRS and Université Paris-Sud.