Only hybrid anyons can exist in broken symmetry phase of nonrelativistic $[U(1)]^2$ Chern-Simons theory

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Abstract

We present two examples of parity-invariant $[U(1)]^2$ Chern-Simons-Higgs models with spontaneously broken symmetry. The models possess topological vortex excitations. It is argued that the smallest possible flux quanta are composites of one quantum of each type (1,1). These hybrid anyons will dominate the statistical properties near the ground state. We analyse their statistical interactions and find out that unlike in the case of Jackiw-Pi solitons there is short range magnetic interaction which can lead to formation of bound states of hybrid anyons. In addition to mutual interactions they possess internal structure which can lead upon quantisation to discrete spectrum of energy levels.
Bogomol’nyi type models in which the ground state will be dominated by topological vortices with winding numbers $(1, 1)$. Vortices of this type are anyons in spite of the fact that the theory is parity invariant. This ground state is a solution with spontaneously broken P-invariance. We also analyse their statistical interactions and come to the conclusion that they can form bound states due to short range magnetic trapping.

1 The model with condensate induced by external magnetic field and normalisability of splitting modes

Let us consider Lagrangian density of the double layer system

$$L = \kappa \varepsilon^{\mu \nu \lambda} a_\mu^{(1)} \partial_\nu a_\lambda^{(2)} + \frac{1}{2} i (\psi^* D_t \psi - \psi D_t \psi^*) - \frac{1}{2} D_k \psi^* D_k \psi$$

$$+ \frac{1}{2} i (\phi^* D_t \phi - \phi D_t \phi^*) - \frac{1}{2} D_k \phi^* D_k \phi - U(\psi, \phi)$$

(1)

where $\kappa$ is chosen to be positive for definiteness. The covariant derivatives and the potential are

$$D_\mu \psi = \partial_\mu \psi - i a_\mu^{(1)} \psi - i A_\mu \psi , \ \ D_\mu \phi = \partial_\mu \phi - i a_\mu^{(2)} \phi - i A_\mu \phi ,$$

$$U = \frac{1}{\kappa} |\psi|^2 |\phi|^2 - \frac{|B_{ext}|^2}{2} (|\psi|^2 + |\phi|^2) .$$

(2)

This is a simplified version of the Lagrangian in [4] and of the nonrelativistic model in [3]. Chemical potentials are proportional to $|B_{ext}|$ and the mixed term of the fourth order can be regarded as a renormalisation counterterm [5, 4]. Coulomb interaction [5, 4] is neglected at the beginning. The chemical potentials can depend on $B_{ext}$ also in some other way, say $\mu(B_{ext})$. This model is Bogomol’nyi type of theory only at such a special value (values) of external magnetic field at which $\mu(B_{ext}) = \frac{1}{2} |B_{ext}|$. We have restricted to the case with purely mutual statistical interaction to provide the paper with greater clarity. The external potentials $A_\mu$ were introduced to provide the uniform external magnetic field background

$$B_{ext} = -\varepsilon_{mn} \partial_m A_n^{ext} , \ \ A_0^{ext} = 0 .$$

(3)

Variation of the Lagrangian with respect to $a_0^{(I)}$, with $I = 1, 2$, leads to Gauss’ laws

$$\kappa B_1 = \phi^* \phi , \ \ \kappa B_2 = \psi^* \psi ,$$

(4)

which should be regarded as constraints with Lagrange multipliers: $a_0^{(I)}$. With such a constraint the Hamiltonian density reads

$$H = \frac{1}{2} D_k \psi^* D_k \psi + \frac{1}{2} D_k \phi^* D_k \phi + U .$$

(5)
Now we apply the Bogomol’nyi trick

\[ D_k \psi^* D_k \psi = |(D_1 + iD_2)\psi|^2 + B_1 \rho_1 + B^{ext} \rho_1 + \nabla \times \vec{J}_1 . \]  

(6)

With this decomposition and with the use of the Gauss’ laws one obtains after neglect of boundary terms, for \( B^{ext} < 0 \)

\[ H = \frac{1}{2} |(D_1 + iD_2)\psi|^2 + \frac{1}{2} |(D_1 + iD_2)\phi|^2 , \]  

(7)

which is positively definite. This Hamiltonian is minimised by fields satisfying following equations

\[ D_+ \psi , \ D_+ \phi = 0 , \]  

(8)

together with Gauss’ laws. Solutions to the above equations are also solutions to the Euler-Lagrange equations of the model provided that Lagrange multipliers are taken equal to

\[ a^{(1)}_0 = \frac{1}{2\kappa} \rho_2 , \ a^{(2)}_0 = \frac{1}{2\kappa} \rho_1 , \]  

(9)

From the self-duality equations (8) we obtain

\[ a^{(I)}_k = \frac{1}{2} \varepsilon_{ki} \partial_l \ln \rho_1 + \partial_k \omega I - A^{ext}_k , \ (1 \leftrightarrow 2) , \ I = 1,2 , \]  

(10)

where \( \omega_{1,2} \) are phases of the fields \( \psi \) and \( \phi \) respectively. Substitution to the Gauss’ laws (8) yields

\[ \nabla^2 \ln \rho_1 = \frac{2}{\kappa} (\rho_2 - \rho_0) + 2\varepsilon_{mn} \partial_m \partial_n \omega_1 , \ (1 \leftrightarrow 2) , \ \rho_0 = \kappa |B^{ext}| . \]  

(11)

Assuming the phases of the Higgs fields to be of the form

\[ \omega_I = \sum_{p_I} \Theta(\vec{x} - \vec{R}_{p_I}) , \ p_I = 1, ..., n_I , \]  

(12)

which fixes the gauge \( \partial_k a^{(I)}_k = 0 \), we get

\[ \nabla^2 \ln \rho_1 = \frac{2}{\kappa} (\rho_2 - \rho_0) + 4\pi \sum_{p_1} \delta^{(2)}(\vec{x} - \vec{R}_{p_1}) , \]  

\[ \nabla^2 \ln \rho_2 = \frac{2}{\kappa} (\rho_1 - \rho_0) + 4\pi \sum_{p_2} \delta^{(2)}(\vec{x} - \vec{R}_{p_2}) . \]  

(13)

In the special case of \( \rho_1 = \rho_2 \equiv \rho \) and outside of the singular points we will obtain

\[ \nabla^2 \ln \rho = \frac{2}{\kappa} (\rho - \rho_0) . \]  

(14)
This equation is known to possess static multivortex solutions \cite{8}. The separate vortices are in fact composites of identical vortices of two different types sitting on top of each another. Natural question arises whether we can expect also existence of separate vortices of different types. Now we will explicitly count the number of normalisable zero modes around the static (1,1) solution. It will appear that there are only two translational modes. The modes which could correspond to splitting of (1,1) vortex into separate (1,0) and (0,1) vortices are not normalisable and thus would lead to unbounded rise of charge with respect to background. We will consider special configuration of coinciding vortices with vorticity \((n,n)\).

Let the unperturbed solution be of the form
\[
\psi = \phi = f(r)e^{in\theta} .
\]
(15)

Now we take perturbations of the phases of the Higgs field to be \(\alpha_1, \alpha_2\) and those of the moduli: \(f_{h_1}, f_{h_2}\). \[
\psi + \delta\psi = f(r)[1 + h_1(r, \theta)]e^{in\theta + i\alpha_1(r, \theta)} ,
\]
\[
\phi + \delta\phi = f(r)[1 + h_2(r, \theta)]e^{in\theta + i\alpha_2(r, \theta)} .
\]
(16)

Linearisation of the self-dual equations \((8)\) with respect to perturbations of the Higgs fields and those of gauge potentials \(c^{(I)}_k\), \(I = 1, 2\), yields
\[
c^{(I)}_{\theta} = \partial_r h_I - \frac{1}{r}\partial_\theta \alpha_I ,
\]
\[
c^{(I)}_r = -\frac{1}{r}\partial_\theta h_I - \partial_r \alpha_I
\]
(17)
for \(I = 1, 2\). Once the perturbations of the Higgs field are known, \(c^{(I)}_k\) can be calculated from the above equations. To have a unique solution we have to fix the gauge
\[
\partial_k c^{(I)}_k \equiv \partial_r c^{(I)}_r + \frac{1}{r}c^{(I)}_r + \frac{1}{r}\partial_\theta c^{(I)}_\theta = 0 \; , \; I = 1, 2 .
\]
(18)

Upon substitution of \((17)\) to this gauge condition we will obtain
\[
\nabla^2 \alpha_1 = 0 \; , \; \nabla^2 \alpha_2 = 0 .
\]
(19)

Similar substitution of \((17)\) to Gauss’ laws \((4)\) will lead to
\[
\nabla^2 h_1 = -\frac{2}{\kappa}f^2(r)h_2 \; , \; \nabla^2 h_2 = -\frac{2}{\kappa}f^2(r)h_1 .
\]
(20)

Now let us consider the following most general Ansatz
\[
h_1(r, \theta) = \sum_{k=1}^n [\xi^1_k P_k(r) \cos k\theta + \xi^2_k Q_k(r) \sin k\theta] ,
\]
\[
h_2(r, \theta) = \sum_{k=1}^n [\lambda^1_k P_k(r) \cos k\theta + \lambda^2_k Q_k(r) \sin k\theta] ,
\]
(21)
with 4n real parameters. This Ansatz leads to following two sets of equations
\[
\xi_1^k[\Delta_k P_k(r)] = [-\frac{2}{\kappa}f^2(r)P_k(r)]\lambda_1^k, \\
\lambda_1^k[\Delta_k P_k(r)] = [-\frac{2}{\kappa}f^2(r)P_k(r)]\xi_1^k, \\
(22)
\]
and
\[
\xi_2^k[\Delta_k Q_k(r)] = [-\frac{2}{\kappa}f^2(r)Q_k(r)]\lambda_2^k, \\
\lambda_2^k[\Delta_k Q_k(r)] = [-\frac{2}{\kappa}f^2(r)Q_k(r)]\xi_2^k, \\
(23)
\]
where the laplacians are
\[
\Delta_k = \frac{d^2}{dr^2} + \frac{1}{r}\frac{d}{dr} - \frac{k^2}{r^2}. \\
(24)
\]
These equations do not lead to contradiction for nonzero \(\lambda, \xi\) only if
\[
\frac{\xi_1^k}{\lambda_1^k} \equiv \sigma_1^k = -1 \quad \text{and} \quad \frac{\xi_2^k}{\lambda_2^k} \equiv \sigma_2^k = +1, \\
(25)
\]
where the choices of signs are independent. Now we have only two equations
\[
P''_k + \frac{1}{r}P'_k + \left[\frac{2\sigma_1^k}{\kappa}f^2 - \frac{k^2}{r^2}\right]P_k = 0, \\
Q''_k + \frac{1}{r}Q'_k + \left[\frac{2\sigma_2^k}{\kappa}f^2 - \frac{k^2}{r^2}\right]Q_k = 0. \\
(26)
\]
Asymptotically at \(r \to \infty\) they become Bessel or modified Bessel equations dependent on whether given \(\sigma_k^i\) is negative or positive, since \(f^2(\infty) = \rho_0\). Thus the only normalisable zero modes are those with \(\sigma_1^k = \sigma_2^k = 1\) or in terms of the parameters
\[
\xi_1^k = \lambda_1^k \quad \text{and} \quad \xi_2^k = \lambda_2^k. \\
(27)
\]
With this condition we have effectively only one equation
\[
P''_k + \frac{1}{r}P'_k + \left[\frac{2}{\kappa}f^2 - \frac{k^2}{r^2}\right]P_k = 0, \quad Q_k = P_k. \\
(28)
\]
The asymptotically vanishing solution possesses singularity at the origin, which can be normalised as
\[
P_k(r) \sim -\frac{1}{r^k}, \quad r \to 0. \\
(29)
\]
With this asymptotics we can choose such a solution to eq.(19) that gauge potentials expressed as in eq.(17) in terms of perturbations of the moduli and
phase of the Higgs field are regular at the origin. This choice appears to be unique
\[
\alpha = \alpha_1 = \alpha_2 = \frac{1}{r^k} \left[ \lambda_1^k \sin k\theta - \lambda_2^k \cos k\theta \right].
\]  
(30)

After we expand the perturbed Higgs fields around the origin we will obtain up to terms linear in \( \lambda \)
\[
\psi + \delta \psi = \phi + \delta \phi \approx \left( z^n - \sum_k \lambda_k z^k \right), \quad \lambda_k = \lambda_1^k + i\lambda_2^k.
\]  
(31)

Now we can clearly see that the only normalisable zero modes are those which correspond to splitting of the coincident \((n, n)\) configuration into a set of hybrid vortices of the type \((p, p)\). Splitting into vortices of the type \((p, q)\) with \(p \neq q\) is forbidden by charge conservation.

We have analysed splitting modes only within the framework of self-dual equations. This analysis shows that the coincident \((1, 1)\) configuration sits at the bottom of potential well. The question now is whether the energy of a pair of \((1, 0)\) and \((0, 1)\) vortices is higher or the same as that of coincident solution. If the latter is the case it would mean that a pair of separate vortices is separated from the coincident configuration by a potential barrier and otherwise degenerate. At low temperatures we would have two coexistent phases of vortex condensates.

We can give a simple argument to rule out the latter possibility. Let us analyse similarly as in \([3]\) asymptotic properties of equations \((13)\). We define small fluctuations around asymptotic values by
\[
g_1 = \frac{2}{\kappa} (\rho_1 - \rho_0), \quad g_2 = \frac{2}{\kappa} (\rho_2 - \rho_0).
\]  
(32)

Equations \((13)\) when linearised with respect to \(g_1\) and \(g_2\) take the form
\[
\nabla^2 g_1 = \alpha g_2, \\
\nabla^2 g_2 = \alpha g_1, \quad \alpha = \frac{2\rho_0}{\kappa} > 0.
\]  
(33)

We can remove \(\alpha\) by rescaling coordinates and then diagonalise this set of equations
\[
\nabla^2 (g_1 + g_2) = \alpha (g_1 + g_2), \\
\nabla^2 (g_1 - g_2) = -\alpha (g_1 - g_2).
\]  
(34)

Nontrivial asymptotic solutions are given by asymptotics of Bessel functions. \((g_1 + g_2)\) vanishes exponentially. An asymptotic decay of \((g_1 - g_2)\) is much slower so only the trivial solution \(g_1 = g_2\) is compatible with finiteness of \(U(1)\) charges. This argument excludes existence of zero energy vortex with winding numbers \((1, 0)\) or \((0, 1)\). Thus the energy of the widely separated pair of \((1, 0)\) and \((0, 1)\) vortices must be higher then the energy of the hybrid anyon with winding numbers \((1, 1)\).

Thus we can expect that at low temperatures the system will be dominated by minimal composites with the winding numbers \((1, 1)\).
2 Model of statistical interaction with uniform background

Another model with dominance of the ground state by hybrid anyons is a direct extension of the U(1) model proposed by Barashenkov and Harin [6].

\[
L = \kappa\varepsilon^{\mu \nu \lambda} a^{(1)}_\mu \partial_\nu a^{(2)}_\lambda + \frac{1}{2} i (\psi^* D_t \psi - \psi D_t \psi^*) - \frac{1}{2} D_k \psi^* D_k \psi
\]

\[
+ \frac{1}{2} i (\phi^* D_t \phi - \phi D_t \phi^*) - \frac{1}{2} D_k \phi^* D_k \phi - V(\phi, \psi) - qa^{(1)}_0 - qa^{(2)}_0 ,
\]

with potential of the form

\[
V(\phi, \psi) = \frac{1}{\kappa} (\rho_1 - q)(\rho_2 - q).
\]

Gauss' laws in this case read

\[
\kappa B_1 = (\rho_2 - q), \quad \kappa B_2 = (\rho_1 - q).
\]

The Hamiltonian after Bogomol'nyi decomposition is

\[
H = \frac{1}{2} |D_+ \psi|^2 + \frac{1}{2} |D_+ \phi|^2 + \frac{1}{2} q (B_1 + B_2).
\]

The energy is bounded from below by

\[
E \geq 2\pi (n_1 + n_2).
\]

This lower bound is saturated by fields satisfying

\[
D_+ \psi = 0, \quad D_+ \phi = 0.
\]

The Lagrange multipliers have to be

\[
a^{(1)}_0 = \frac{1}{2\kappa} (\rho_2 - q), \quad a^{(2)}_0 = \frac{1}{2\kappa} (\rho_1 - q).
\]

Further analysis similar as in section 1 shows that also in this case vortex with winding numbers \((n, n)\) can decay only into vortices of the type \((p, p)\) and not in those of the type \((p, q)\). Once again we are lead to conclusion that hybrid anyons will dominate statistically at low temperatures.

3 Statistics and magnetic interactions of hybrid anyons

Now we will assume that the ground state of the theory is dominated by elementary hybrid anyons of the type \((1, 1)\). We will apply Manton's approximation
to investigate their statistical properties. Namely we assume that a good approximation to configuration of slowly moving vortices are static multivortex solutions with their parameters promoted to the role of time dependent collective coordinates. This method was first applied in the context of relativistic Chern-Simons vortices by Kim and Min \[9\]. In our Lagrangians \[1\] or \[35\] we have first order time derivatives so this prescription can be directly applied only to obtain those terms of the effective Lagrangian which are linear in time derivatives of the parameters. The part of the effective Lagrangian linear in velocities reads

\[ L_{\text{eff}}^{(1)} = \int d^2 x \left[ -\kappa \varepsilon^{mn} a^{(1)}_m \partial_t a^{(2)}_n - \rho_1 \partial_t \omega_1 - \rho_2 \partial_t \omega_2 \right] , \quad (42) \]

The first term gives no contribution since the gauge potentials satisfy the Coulomb gauge \( \partial_k a_k^{(1)} = 0 \) at every moment of time. Now we restrict to the dominant multivortex solution, which satisfies \( \psi = \phi, a^{(1)}_{\mu} = a^{(2)}_{\mu} \),

\[ L_{\text{eff}}^{(1)} = -2 \int d^2 x \left( \rho \partial_t \omega \right) = -2 \kappa \int d^2 x B \partial_t \omega , \quad (43) \]

where we have made use of the Gauss’ laws \[8\] or \[38\]. With the form of the phase

\[ \dot{\omega} = \frac{d}{dt} \sum_p \Theta(\vec{x} - \vec{R}_p) = \sum_p \varepsilon_{ij} \vec{R}_p^i \partial_j \ln | \vec{x} - \vec{R}_p | \]

and some integration by parts we get the form of effective Lagrangian useful for further discussion.

\[ L_{\text{eff}}^{(1)} = -4 \pi \kappa \sum_p \vec{R}_p^i \varepsilon_{ij} a_j(\vec{R}_p) , \quad (45) \]

which is expressed through the gauge field at the cores of vortices and their velocities. Now we can easily get the orbital part of the angular momentum at vanishing velocities

\[ J_{\text{orb}} = \sum \varepsilon_{ij} \vec{R}_p^i \frac{\partial L_{\text{eff}}^{(1)}}{\partial R_p^j} = -4 \pi \kappa \sum_p R_p^i \varepsilon_{ij} a_j(\vec{R}_p) . \quad (46) \]

It is worth noticing that the effective Lagrangian \[43\] takes the form of the coupling of the point particle current to external gauge field defined on the moduli space. Due to this coupling vortex at \( \vec{R}_p \) feels the magnetic field

\[ B_{\text{eff}}(\vec{R}_p) = -4 \pi \kappa \varepsilon_{ij} \partial_i a_j(\vec{R}_p) , \quad \partial_i \equiv \frac{\partial}{\partial R_p^i} . \quad (47) \]

Now let us consider configuration of two vortices each of the type \( (1, 1) \) in the center of mass frame at actual positions \( +\vec{R} \) and \( -\vec{R} \). The orbital momentum reads

\[ J_{\text{orb}} = -4 \pi \kappa [ R^i \varepsilon_{ij} v_j(\vec{R}) - R^i \varepsilon_{ij} v_j(-\vec{R})] = -8 \pi \kappa R v_9(\vec{R}) , \quad (48) \]
where the last equality holds due to the symmetry of the configuration. It is easy to see that

$$ \frac{1}{R} \frac{dJ(R)}{dR} = -8\pi\kappa \frac{v_\theta}{R} + \frac{dv_\theta(R)}{dR} = B_{eff}(+\vec{R}) + B_{eff}(-\vec{R}) = B_{eff}(R). \quad (49) $$

Thus if the spin dependence of two-vortex static configuration on their distance were known the magnetic interaction could be extracted due to the above formula. The formula makes use of the global properties of the solution so it is not very sensitive to numerical errors. The total angular momentum $J$ may differ from its orbital part only by an additive normalisation constant.

Explicite calculation of the linear part of effective Lagrangian from eq. (43) leads for well separated vortices to

$$ L^{(1)}_{eff} = 4\pi\kappa \sum_{p>q} \hat{\Theta}(\vec{R}_p - \vec{R}_q), \quad (50) $$

Short range interactions are more complicated. We can promote $\lambda$-s in eq. (21) in the case $n = 2$, $k = 2$ to the role of collective coordinates. The effective Lagrangian up to leading terms in $\lambda$ can be evaluated as

$$ L^{(1)}_{eff} = b_{eff} \Lambda^2 \omega, \quad b_{eff} = \pi \int_0^\infty r dr \frac{f^2 P_2}{r^2}, \quad (51) $$

where $\lambda^2_1 + i\lambda^2_2 = \Lambda e^{i\omega}$. Positions of vortices are complex square roots of $\lambda$. If they are $+\vec{R}$ and $-\vec{R}$ then in polar coordinates $R^2 e^{2i\Theta} = \Lambda e^{i\omega}$ and

$$ L^{(1)}_{eff} = 2b_{eff} R^4 \hat{\Theta}, \quad (52) $$

which describes the short range asymptotics of magnetic interaction.

So far we have considered only statistical and magnetic interactions of hybrid anyons as a whole. The hybrid anyon can be thought of as a pair of $(1,0)$ and $(0,1)$ vortices trapped in a potential well. In [11] it was shown that there are mutual magnetic short range interactions between the different types of vortices. In the center of mass frame the internal dynamics of single hybrid anyon reduces to that of short range oscillator in external magnetic field. Quantum mechanically we can expect an interesting discrete spectrum.

### 4 Conclusions

In both cases analysed in this paper we are lead to the conclusion that the dominant flux quanta are vortices with winding numbers $(1,1)$. In the second model they also have smallest energy thanks to the Bogomol’nyi lower bound which they saturate and thus will dominate statistically. In the case of the first model all solutions are degenerate in energy. Degeneracy can be removed by
introduction of Coulomb interaction which will raise the energy of excitations with bigger winding number equivalent to bigger charges.

The short range remnant of statistical interaction is rather a magnetic interaction unlike in the case of Jackiw-Pi solitons investigated by the same method [13]. Thus hybrid anyons can form bound states due to the possibility of magnetic trapping [10, 11]. At the quantum level the effect of these classical trapped states can be observable in the form of resonances characterised by Landau levels. This also means that when one constructs effective field theory for vortices like in [10, 11] the counterterm due to the short range statistical interaction can be much smaller than expected. Due to the finite width of topological vortices the statistical interaction is effectively switched off at short distances. Also the internal structure of the hybrid anyon can manifest itself by excitations with discrete energy levels.

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