Deconfinement of vortices with continuously variable fractions of the unit flux quanta in two-gap superconductors

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Abstract – We propose a new stage of confinement-deconfinement transition, which can be observed in laboratory. In two-gap superconductors (SCs), two kinds of vortex exist and each of them carries a continuously variable fraction of the unit flux quanta $\Phi_0 = hc/2e$. The confined state of these two is a usual vortex and stable in the low-temperature region of the system under a certain magnetic field above $H_{c1}$. We see an analogy to quarks in a charged pion. An entropy gain causes two fractional vortices to be deconfined above a certain temperature.

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Introduction. – Topological defects furnish fascinating problems in physics, and keep attracting a lot of concerns. In general, they lead the quantization of physical quantities. Vortex is one of the most important example of topological defects [1]. It is well known that in superconductors (SCs), the requirement that the order parameter is a single-valued function guarantees the quantization of the magnetic flux of a vortex in the unit of $\Phi_0 = hc/2e$.

The fractionalization of topologically quantized numbers have also been researched extensively in many different contexts [2]. Vortices with fractional fluxes have been discussed in $p$-wave superfluid $^3$He [3] and $d$-wave SCs with a grain boundary [4] (see also a notice1). The fractionalization comes from unconventional structures of order parameters for these superfluid and SCs without any contradiction to their single-valued nature.

Here we discuss the vortex state in two-gap SCs [5] which has two superconducting gaps on two separate Fermi surfaces. The two-gap superconductivity can be seen in a lot of SCs in solid state and has attracted special attention. MgB$_2$ is a recent typical textbook [6]. The possibility of the two-gap superconducting state is also pointed out in 2H-NbSe$_2$ [7], PrOs$_4$Sb$_{12}$ [8], $Y_2C_3$ [9], Ca-doped YBCO [10], and liquid hydrogen [11]. The fractionalization of the flux quanta has also been discussed in two-gap SCs [12,13]. Tanaka considered a wire ring and showed that a trapped flux in the ring can be an arbitrary fraction of $\Phi_0$ [12]. Babaev considered a planer geometry and discussed the spontaneous fractional vortices associated with the Berezinskii-Kosterlitz-Thouless (BKT) transition [14] in the weak-coupling limit of an internal Josephson interaction with respect to the relative phase of the two gap functions [13].

Besides these pioneering works, searching for the appearance of the fractional vortices under other setups is challenging and important since it gives new viewpoints of the problem and also new possibility for experimental discovery. In this paper, we propose a novel mechanism for the appearance of fractional vortices in the two-gap SCs. We consider the system under an applied magnetic field and show that a vortex with the unit flux is divided into two vortices with continuously variable fraction of $\Phi_0$ by entropic effect. This phenomena can be seen as deconfinement of two fractional vortices. We see that in the confined state, two vortices behave like quarks in the charged pion [14]. Hence the present theory gives rise to a new stage of confinement-deconfinement transition, which can be examined in laboratory.

Fractional vortices in a two-gap SC. – Let us discuss the property of fractional vortices in a type-II two-gap SC. We consider two gap functions $\Delta_L$ and $\Delta_S$ $|\Delta_L| \gg |\Delta_S|$ opened on two different Fermi surfaces. Their pairing symmetry does not matter. We can consider not only $s$-wave but also unconventional pairing states [15]. We apply a magnetic field along the $z$-direction to a sample of thin-film geometry, i.e., $d < \min(\xi_L^z, \xi_S^z)$, where $d$ is the sample thickness and $\xi_L^z$ and $\xi_S^z$ are

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1Present affiliation: (C) HONDA Motor Co., Ltd.
2The fractional vortex is defined as a vortex with a flux less than $\Phi_0/2$ in ref. [4], while less than $\Phi_0$ in this paper.
the coherence lengths of the gap functions along the z-direction. The gap functions are then expressed as, \( \Delta_{L,S}(r) = |\Delta_{L,S}(r)| \exp(-i \theta_{L,S}(r)) \), where \( r \) is the two-dimensional vector in the \( xy \)-plane. The Ginzburg-Landau (GL) free energy for a two-gap system is the straight extension of the conventional one and given in ref. [16]. We consider the GL free energy in the London limit, namely, we neglect the spacial dependence of \( |\Delta_{L}(r)| \) and \( |\Delta_{S}(r)| \) and obtain

\[
F_{\text{GL}} = d \int d^2r \left[ \frac{\mathbf{H}^2}{8\pi} + \frac{1}{8\pi \lambda_L^2(T)} \left( \mathbf{A} - \frac{\Phi_0}{2\pi} \nabla \theta_L \right)^2 + \frac{1}{8\pi \lambda_S^2(T)} \left( \mathbf{A} - \frac{\Phi_0}{2\pi} \nabla \theta_S \right)^2 - \Gamma(T) \cos(\theta_L - \theta_S) \right].
\]

The characteristic points of this free energy are the presence of two typical length scales, \( \lambda_{L,S}(T) = (4e^2 K_{L,S}|\Delta_{L,S}(T)|^2/\hbar c^2)^{-1/2} \), where \( K_{L,S} \) are the coefficients of the gradient terms in the GL free energy, and of the Josephson-type interaction with a temperature-dependent coupling \( \Gamma(T) \equiv 2\pi/|\Delta_L(T)||\Delta_S(T)| \). This coupling causes the locking effect for the relative phase \( \theta_L - \theta_S \). The equation of motion for the gauge field \( \mathbf{A} \) is obtained from \( \delta F_{\text{GL}}/\delta \mathbf{A} = 0 \) as,

\[
\lambda(T)^2 \nabla \times \mathbf{H} = -\mathbf{A} + \frac{\Phi_L(T)}{2\pi} \nabla \theta_L + \frac{\Phi_S(T)}{2\pi} \nabla \theta_S,
\]

where

\[
\lambda^2(T) \equiv \frac{1}{\lambda_L^2(T)} + \frac{1}{\lambda_S^2(T)}
\]

is the London penetration depth, and

\[
\Phi_L(T) + \Phi_S(T) = \Phi_0.
\]

The following boundary conditions are imposed:

\[
\nabla \times \nabla \theta_{L,S}(r) = 2\pi \hat{z} \delta^2(r - \mathbf{R}_{L,S}),
\]

which would be relevant for the dilute vortex system with \( H \approx H_{c1} \). The singlevaluedness of two gap functions is guaranteed except for \( \mathbf{R}_L \) and \( \mathbf{R}_S \) because of the 2\pi winding. From eq. (2), we obtain the London equation for the two-gap case:

\[
\lambda^2(T) \nabla \times \nabla \times \mathbf{H} + \mathbf{H} = \Phi_L(T) \hat{z} \delta^2(r - \mathbf{R}_L) + \Phi_S(T) \hat{z} \delta^2(r - \mathbf{R}_S).
\]

This equation tells us important implications. When \( \mathbf{R}_L = \mathbf{R}_S = \mathbf{R} \), the equation reduces to the usual London equation \( \lambda^2(T) \nabla \times \nabla \times \mathbf{H} + \mathbf{H} = \Phi_0 \hat{z} \delta^2(r - \mathbf{R}) \), and we have a vortex with the unit flux passing through \( \mathbf{R} \). When \( \mathbf{R}_L \neq \mathbf{R}_S \), however, we have two vortices with \( \Phi_L(T) \) through \( \mathbf{R}_L \) and \( \Phi_S(T) \) through \( \mathbf{R}_S \). It should be emphasized that \( \Phi_L(T) \) and \( \Phi_S(T) \) are fractions of \( \Phi_0 \) and its ratio is determined by the ratio \( \lambda_L(T)/\lambda_S(T) \) (see, eqs. (3), (4), and (5)). When \( \lambda_L(T)/\lambda_S(T) \) is limited to unity (such a limitation could come from the crystal symmetry, if we consider unconventional states in two-dimensional irreducible representations in a system with single Fermi surface [15]), we have only the half-flux \( \Phi_L(T) = \Phi_S(T) = \Phi_0/2 \), but there is no such a limitation in the present system and fractional vortices besides the half-flux vortex can exist when \( \lambda_L(T) \neq \lambda_S(T) \). It should be also emphasized that these fractional fluxes vary continuously with \( |\Delta_L(T)|/|\Delta_S(T)| \), which determines the ratio \( \lambda_L(T)/\lambda_S(T) \).

There is a possibility to observe these fractional vortices when their separation is large enough. It is small in the low-temperature region due to the Josephson-type interaction in the GL free energy (1), but is possible to be large in the high-temperature regime in order to gain entropy. To check this scenario, we calculate the energy cost and the entropy gain for a certain configuration of two vortices with \( \Phi_L(T) \) and \( \Phi_S(T) \).

**Configuration energy.** — We can eliminate the gauge field from eq. (1) by using the equation of motion (2), and obtain the energy cost,

\[
H_C = d \int d^2r \left[ \frac{\mathbf{H}^2}{8\pi} + \frac{\lambda^2(T)}{8\pi} (\nabla \times \mathbf{H})^2 + \frac{K(T)}{2} \left( \nabla (\theta_L - \theta_S) \right)^2 - \Gamma(T) \cos(\theta_L - \theta_S) \right],
\]

where \( K(T) \equiv \{ \Phi_L(T)\Phi_S(T) \}/16\pi^3 \lambda(T)^2 \). The vortex solution in the thin-film geometry for the single gap SC has been investigated [17]. The extension to the two-gap case is straightforward. We introduce the effective London penetration depth in the thin film \( \lambda_{eff} = \lambda_L^2/\lambda_S(T) \). When \( \lambda_{eff} \) is comparable to the system size, we obtain

\[
H_C \approx \sum_{i=L,S} \epsilon_i(T) - 2\pi K(T) \ln \frac{\mathbf{R}_L - \mathbf{R}_S}{\lambda_{eff}(T)} + E_{rel},
\]

where

\[
\epsilon_{L,S}(T) \equiv \{ \Phi_{L,S}^2(T) / 16\pi^2 \lambda_{eff}(T) \} \ln(\lambda_{eff}(T)/\xi_{L,S}(T))
\]

are the creation energy for L- and S-vortices, \( \xi_{L,S}(T) \equiv \hbar v_F / |\Delta_{L,S}(T)| \) the coherence lengths of L- and S-gap functions, each of which should be in the same order of each vortex core, and \( v_F \) the Fermi velocities. The logarithmic term on the r.h.s is the magnetic interaction between two vortices, which is repulsive and shows a logarithmic behavior. The last term has a sine-Gordon form

\[
E_{rel} = d \int d^2r \left[ \frac{K(T)}{2} \left( \nabla (\theta_L - \theta_S) \right)^2 - \Gamma(T) \cos(\theta_L - \theta_S) \right].
\]
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When $\gamma = 0$, $\Gamma(T) = 0$ and $E_{rel} \propto \ln |R_L - R_S|$ [13]. Such a case can be applied to superconductivity in liquid hydrogen [11]. We go beyond the weak-coupling limit of $\gamma$, which would be significant for the two-gap superconductivity in solid state, e.g., MgB$_2$ [6].

In eq. (10), the spatial dependence of the relative phase loses the elastic energy. The Josephson-type term favors $\theta_L - \theta_S = 2\pi n$ for $\gamma > 0$ and $\theta_L - \theta_S = (2n + 1)\pi$ for $\gamma < 0$. But we have the boundary condition (6). There should be a line-like singularity between separated vortices associated with a 2$\pi$ jump of the relative phase. This “string” is the two-dimensional extension of the kink in the sine-Gordon model in one dimension [18]. Therefore, for large $\gamma$,

$$E_{rel} = \epsilon_{st}(T)L_{st} + \text{const}, \quad (11)$$

where $\epsilon_{st}(T) \equiv d\sqrt{K(T)\Gamma(T)}|/\pi + O(1/\gamma)$ is the string creation energy per a unit length and $L_{st} \gg |R_L - R_S|$ is the length of the string. To minimize $H_{c_1}$, $L_{st} \rightarrow |R_L - R_S|$. This indicates that there is a linear confinement potential between L- and S-vortices. We have also the two-dimensional Coulomb type repulsion. Then, two fractional vortices are analogous to massive “quarks” with “fractional electric charge” $\Phi_L$ and $\Phi_S$, except for the spatial dimensionality [14]. The L-S confined state corresponds to the “charged pion” (the total charge +1 of the pion corresponds to the unit flux $\Phi_0 = \Phi_L + \Phi_S$ of the L-S confined state).

Entropy, confinement, and deconfinement. – Let us examine the thermodynamic stability of the separated state with a distance $R \equiv |R_L - R_S| = N\xi \neq 0$, where $\xi \equiv \max(\xi_L, \xi_S)$ and we refer to it as the cut-off length of this system at short length scale. We introduce an artificial square lattice with a constant $\xi$ so that two vortices sit on one of the lattice axes. We discuss the free energy $F = H_{c_1} - TS$, where $H_{c_1}$ is given in eqs. (9) and (11). There are two contributions to the entropy $S = S_{loc} + S_{st}$, the former one comes from the location of two vortex cores $S_{loc} \simeq 2k_B\ln(\Omega/\xi^2)$ ($\Omega$ the area of the system) and the latter the string configuration $S_{st}$.

The lowest-energy configuration of the string is the straight line connecting two vortices, where $L_{st} = R$. The kink-like configurations of the string would be thermally excited and each excitations contributes $\xi$ to $L_{st}$ (see fig. 1). As shown in fig. 1, there are two kinds of excitations one of which is directed “upward” and the other “downward”. The numbers of the upward and downward excitations are equal to each other so that the string is ended by two fixed points (vortices) and we denote them as $n$.

We consider states with $L_{st} \simeq R$, i.e., $n \ll N$. By using the dilute approximation, we find $S_{st} \simeq k_B\ln c_{N-n}c_n \simeq k_B(\delta R/\xi)\ln(\delta R/2R)$, where $\delta R = 2n\xi$. We find that the value of $\delta R$ which minimizes $F$ at fixed $R$ is $\delta R_0(R) = 2R\exp[-\xi_{st}/k_BT]$. Let us define a function $\tilde{F}(R) \equiv F(R, \delta R_0(R))$. We introduce $\alpha(T)$ as the value of the vortex distance $R$ at the minimum of $\tilde{F}(R)$ and that shows a diameter of the confined region of two fractional vortices. We obtain

$$a(T) = \frac{2\pi K(T)/\epsilon_{st}(T)}{1 - \frac{2k_BT^2}{\epsilon_{st}(T)} \exp \left[-\frac{\xi_{st}k_BT}{2k_BT}\right]}. \quad (12)$$

Remind that the solution of $1 - 2/x \cdot \exp[-x] = 0$ is $x = -\text{LambertW}(2) \simeq 0.85$. Then, $a(T) \rightarrow \alpha$ (see footnote 2) when $k_BT = \frac{\xi(T)\epsilon_{st}(T)}{\text{LambertW}(2)} = 0$. \(13\)

The temperature that satisfies eq. (13) would correspond to the deconfinement temperature and we refer to it as $T_{dec}$. We emphasize that it must surely be $T_{dec} < T_c$, where $T_c$ is the gap opening temperature. The reason is that $\xi(T)\epsilon_{st}(T)$ monotonically decreases with increasing $T$ and vanishes at $T_c$. Then, the l.h.s. of eq. (13) changes its sign once between zero temperature and $T_c$. This statement is crucial since vortices cannot exist above $T_c$ and the argument becomes inconsistent when $T_{dec} > T_c$.

We can also show the presence of $T_{dec}$ when we consider states with $L_{st} \gg R$. Let us remind that the number of configurations $N_{con} = \mu^n$ for the n-step self-avoiding random walk, where $\mu$ is a constant and equal to 3 for the 2D square lattice case. For a string with sufficiently large $L_{st}$, $N_{con}$ would be well approximated by the random walk result and we obtain $N_{con} \simeq \mu^{L_{st}/\xi}$. Then, the entropy of the string is proportional to $L_{st}$ [14] and the string free energy $F_{st} \simeq \tau_{st}(T)L_{st}$, where the coefficient $\tau_{st}(T) \equiv \{\epsilon_{st}(T) - k_BT/\xi(T) \cdot \ln \mu\}$ denotes the string tension and is positive in the low-temperature region, i.e., fractional vortices are confined. Above a certain temperature, $\tau_{st}(T) < 0$, namely, the string loses its tension and deconfinement occurs. $T_{dec}$ is obtained from the equation $\tau_{st}(T_{dec}) = 0$, which is equivalent to eq. (13) when we take $\mu = \exp[\text{LambertW}(2)] \simeq 2.35$.

To obtain $T_{dec}$, we should know the temperature dependence of the two gaps. By using the GL equation for the two-gap superconductor with sufficiently large $\gamma$ [16], we check that $|\Delta_{LS}(T)|/|\Delta_{LS}(0)| = k\sqrt{T/T_c}$, where $t \equiv T/T_c$ and $k$ is a constant which can be obtained

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2Actually, $\alpha(T)$ would have a cut off determined by the inverse of the square root of the applied magnetic field.
Fig. 2: $A$ dependence of $t_{\text{dec}} \equiv T_{\text{dec}}/T_c$, where $T_{\text{dec}}$ is the deconfinement temperature.

Table 1: Parameters which would be relevant for MgB$_2$ [6].

| Parameter | Value |
|-----------|-------|
| $\xi(0)/\lambda(0)$ | 0.32 |
| $d$ | 10 Å |
| $\alpha$ | 2.2 |
| $k$ | 2.1 |

We numerically. Then, eq. (13) becomes $A\sqrt{1-t/t} = \text{LambertW}(2)$, where, in the case of $K_S/K_L = 1$ [16],

$$A = \frac{\Phi_0}{4\pi^{3/2}} \left\{ \xi(0)/\lambda(0) \right\} \frac{2\gamma x}{(x+1/x)^2}$$

and $x \equiv |\Delta_S(0)/|\Delta_L(0)|$. We see from fig. 2 that $t_{\text{dec}} \equiv T_{\text{dec}}/T_c$ depends on $A$ strongly when $A \sim O(1)$. By using the parameters in table 1 which would be relevant for MgB$_2$, we obtain $A \sim 5.7$ and $t_{\text{dec}} \sim 0.97$.

We discuss the BKT transition temperature associated with the usual vortex ($\theta_L = \theta_S$) obtained from $k_S T_{\text{BKT}} = \Phi_0^2/32\pi^2 \lambda_{c,f}(T_{\text{BKT}})$ (see eq. (1)). By using the parameters in table 1, we see $t_{\text{BKT}} \sim 0.90$, which is lower than the $t_{\text{dec}}$ we have obtained. The present argument for $t_{\text{dec}}$ is relevant for the dilute vortex system but does not include the multi-vortex effect which would be important above $T_{\text{BKT}}$, since a lot of thermally activated vortices and anti-vortices exist above $T_{\text{BKT}}$ even in the system under the applied magnetic field considered here [19]. Then, the result would be modified in this case. It should be noted that the coefficient $A$ includes $\xi(0)$ explicitly, on the other hand, the equation for $T_{\text{BKT}}$ does not. So, $T_{\text{dec}}$ can be lower than $T_{\text{BKT}}$ in a system with sufficiently small $\xi(0)$ and the present argument can be precise. It is well known that high-$T_c$ cuprates have rather small coherence length in the order of 10 Å. Therefore a two-gap superconductor Ca-doped YBCO [10] seems to be one of the best candidates to apply the present argument.

We comment on the relation to the renowned decoupling transition [20]. The renormalization group (RG) approach of the 2D sine-Gordon model discussed in ref. [20] suggests that the cosine term in eq. (10) becomes irrelevant above a certain temperature $T_*$, which indicates that decoupling between $\theta_L$ and $\theta_S$ occurs. The lower bound of $T_*$ is obtained from $T_* = 8\pi K(T_*)d$ in the limit $\gamma \to 0$ [20]. This RG analysis takes into account string loops (closed strings) only, i.e., it is assumed that the relative phase $\theta_L - \theta_S$ is single valued. The open string (see fig. 1) should be taken into account since it would be created by the thermal fluctuations. Actually, we have shown that the open string is stabilized by the entropic effect and that deconfinement occurs at $T_{\text{dec}}$. For sufficiently small $A$, $T_{\text{dec}}$ can be lower than the lower bound value of $T_*$, i.e., deconfinement can occur even in the temperature region where the cosine term is relevant.

We also find that

$$\frac{a(T)}{a(0)} = \left[ 1 - \frac{2t}{A\sqrt{1-t/t}} \exp \left\{ -A\sqrt{1-t/t} \right\} \right]^{-1}, \quad (14)$$

where $a(0) = 2\pi K(0)/\epsilon \alpha(0)$. The plot of $a(T)$ is given in fig. 3 for $A = 1.6$. We see that even in the confined phase it seems possible to observe each of the fractional vortices in the region of large $a(T)$ by using some appropriate vortex imaging techniques.

Fig. 3: $a(T)$ for $A = 1.6$. In this case, $t_{\text{dec}} \sim 0.81$. Note that the unit of the vertical line is $\xi(T) = \xi(0)/\sqrt{1-t}$. Here, we use the relation $a(0) \equiv 0.3\xi(0)$ obtained from table 1.

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