Topological defect densities in type-I superconducting phase transitions

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We examine the consequences of a cubic term added to the mean-field potential of Ginzburg-Landau theory to describe first-order superconducting phase transitions. Constraints on its existence are obtained from experiment, which are used to assess its impact on topological defect creation. We find no fundamental changes in either the Kibble-Zurek or Hindmarsh-Rajantie predictions.

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I. INTRODUCTION

It is generally believed that the Universe, evolving from the initial Big Bang, underwent a series of symmetry-breaking phase transitions[1,2] accompanied by the creation of topological defects, frustrations of the unbroken phase within the broken one, induced by continuity of the order-parameter values. These defects appear as magnetic monopoles, cosmic strings, domain walls, and textures.

Direct experimental tests of these ideas are unfeasible, but transitions described by similar equations occur in experimentally accessible condensed-matter systems, and a new trend has unfolded which compares these two systems. This "cosmology in the laboratory" relies on the fact that the dynamics of phase transitions lie in universality classes and the cosmological ones are hence analogous to those of condensed matter. For instance, vortices created in the superfluid phase transitions of ^{4}He and ^{3}He have been studied experimentally (see, e.g., Refs. 3 and 4 for extensive discussions) following an earlier suggestion in which common features with cosmic strings have been noted.[5] Similarities between cosmological phase transitions and the isotropic-nematic phase transition in liquid crystals were studied in Refs. 6 and 7. Analogy with the thermodynamics and transitions in polymer chains was drawn in Ref. 8.

The case of superconductors is of particular interest as the associated phase transition involves a local gauge symmetry-breaking process. In superconductors, cosmic strings manifest themselves as flux tubes or vortices. Two experiments aimed at observing defect densities in superconductive transitions have led to conflicting results.[9,10] In the first, a thin sample of $YBa_2Cu_3O_{7-\delta}$ was laser-heated above T_c, then allowed to relax back to a superconducting state under the cooling power of the refrigerator; a surrounding superconducting quantum interference device (SQUID) loop with a sensitivity of 20 Φ_0/cm^2 was used to detect the net flux creation. The experiment yielded negative results at the level of 10^{-3} of the Kibble-Zurek (KZ) mechanism.[9] Explanations for this include an overly optimistic estimate of the vortex-antivortex difference, the lack of validity of the prediction for a local symmetry breaking, and dissipation of the flux structure more rapidly than experimentally detectable. In contrast, the second experiment, involving a series of ~ 200 connected Josephson junctions under the same experimental conditions, yielded a net flux creation about twice the KZ prediction.[10] The KZ mechanism is, however, accurate only for global symmetry breaking, a situation where the geodesic rule for phase angle summation is valid. A local gauge treatment by Hindmarsh-Rajantie (HR) (Ref. 11) identifies a new mechanism for defect generation, which leads to a prediction for the first experiment well below the measurement sensitivity, although the prediction for the second experiment is in reasonable agreement with observation.

The above experiments were both conducted in type-II materials, which exhibit a second-order phase transition. The question naturally arises as to the extent of changes in the defect density predictions for type-I superconductors. This is the motivation for our work.

Distinction between type-I and type-II superconductors is traditionally made through the Ginzburg-Landau (GL) parameter $\kappa=\lambda/\xi$, the ratio between the magnetic-field penetration length, λ, and the order-parameter (scalar field) coherence length, ξ. These characteristic length scales are obtained, in the presence of a gauge field $\vec{A}$, from the free energy density

$$F(\Phi)=\frac{1}{2m_e}\left|i\hbar\vec{\nabla}\Phi-\frac{e}{c}\vec{A}\Phi\right|^2+V(\Phi)+\frac{1}{2}\vec{\mu}\cdot(\vec{\nabla}\times\vec{A}), \tag{1}$$

where $\vec{\mu}$ is the magnetic moment of the specimen, m_e is the electron mass, and Φ is the order parameter. The GL potential is usually written as[5]

$$V(\Phi)=\alpha\Phi^2+\frac{\beta}{2}\Phi^4, \tag{2}$$

where α is assumed to depend linearly on the temperature, $\alpha=\alpha'(t-1)$, $t\equiv T/T_c$, α' and β are constants, and T_c is the critical temperature. Thus one obtains

$$\lambda=\sqrt{\frac{m_e c^2}{4\pi e^2}\frac{\beta}{|\alpha|}} \tag{3}$$

and

$$\xi=\frac{\hbar}{\sqrt{2m_e|\alpha|}}. \tag{4}$$

At $T=0$ the coherence length is given by $\xi_0=\hbar/\sqrt{2m_e\alpha'}$, with $\kappa\sim\sqrt{\beta}$. For $\kappa>1/\sqrt{2}$, the transition is second order, and ξ_0 is typically less than $\sim 0.04\ \mu$m; for $\kappa<1/\sqrt{2}$, the transition is first-order, with ξ_0 typically greater than $\sim 0.08\ \mu$m. In general, both are second order for $H=0$. First-order transitions arise from the external field term in Eq. (1) in the event that a characteristic sample dimension is greater than λ.

In thermal field theory (TFT) a first-order phase transition arises from consideration of one-loop radiative corrections to a potential of the form of Eq. (2), which introduce a barrier between minima through a cubic scalar field term, as

$$V(\Phi)=\alpha\Phi^2-\gamma|\Phi|^3+\frac{\beta}{2}\Phi^4, \tag{5}$$

where $\gamma(T)=(\sqrt{2}/4\pi)e^3T$.[12] As in the previous case, β is constant and $\alpha=\alpha'(t-1)$ is linear with temperature.

A similar $-\gamma|\Phi|^3$ term, however, arises in considerations of gauge-field fluctuations in the normal-to-superconductor phase transition,[13,14] with

$$\gamma=8\mu_0\frac{e}{\hbar c}\sqrt{\pi\mu_0}T_c, \tag{6}$$

resulting in a first-order phase transition for all values of κ. Crossovers between first- and second-order transitions arise in considerations of thermal fluctuations,[14] as also in nonlocal BCS treatments.[15]

With this is mind, we adopt a potential of the form of Eq. (5) and explore constraints on $\gamma(T)$ from experiment. The results are compared with TFT one-loop radiative corrections. Comparison with the results of Ref. 13, which deals with temperatures close to T_c, is attained in the limit $t\to 1$.

These results are then used to analyze the impact of the cubic term on the KZ and the HR predictions for type-I superconductors, namely, on the models themselves and on a possible nucleation suppression due to the slowing down of the transition induced by the potential barrier.

II. TEMPERATURE SENSITIVITY BOUND

A first-order superconducting phase transition manifests supercritical fields under variation of the temperature, as shown in Fig. 1. The superheating curve is given by the condition $dV/d\Phi=d^2V/d\Phi^2=0$, for a certain value of $\Phi\neq 0$. This yelds $\alpha=9\gamma^2/16\beta$.

In contrast, the supercooling curve is given by the condition $d^2V(\Phi)/d\Phi^2=0$ for $\Phi=0$, corresponding to $\alpha=0$. The unobservable critical curve is given by the condition $V(0)=V(\Phi_c)$ and $dV(\Phi_c)/d\Phi=0$, where Φ_c is the nonvanishing minimum of the potential. This corresponds to $\alpha=\gamma^2/2\beta$.

Introducing the linear dependencies $\alpha=\alpha'(t-1)$ and $\gamma(t)=\delta t$, we obtain for the superheating curve

$$\alpha'(t-1)=\frac{9}{16}\frac{\delta^2}{\beta}t^2, \tag{7}$$

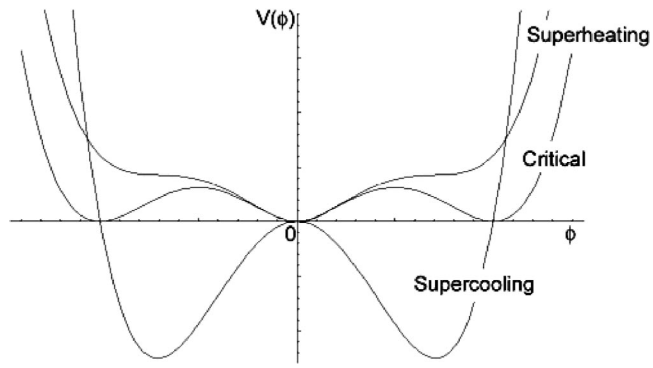

FIG. 1. Characteristic potential curves.

and hence

$$t_{sh}=\frac{2}{1+\sqrt{1-\frac{9}{4}\frac{\delta^2}{\alpha'\beta}}}\sim 1+\frac{9}{16}\frac{\delta^2}{\alpha'\beta}. \tag{8}$$

The superheating curve converges to a zero-field shift in temperature from T_c by $(9\delta^2/16\alpha'\beta)T_c$. This shift is undetected at the current experimental temperature sensitivity of $\Delta t_{exp}\sim 10^{-3}$.[16] Therefore, a bound for the slope of γ is

$$\frac{9}{16}\frac{\delta^2}{\alpha'\beta}<\Delta t_{exp}. \tag{9}$$

The supercooling transition occurs at $t=1$, the critical temperature, since it is determined solely by $\alpha=0$ (this neglects a small correction to α, as given by Ref. 13).

III. SUPERHEATING PERTURBATION BOUND

Measurements of the supercritical fields have commonly been performed on microspheres of type-I materials as a means of determining κ. Table I indicates the critical properties of, Sn and Al. The existence of a cubic term should also be manifest in the presence of a magnetic field. To assess its influence, we repeat Ref. 23 calculations, including the cubic term in the potential. For a small superconducting sphere of radius a, the magnetic moment is given by[23]

$$\frac{\mu}{V}=-3\left[1-\frac{3\lambda}{a\Phi_0}\coth\frac{a\Phi_0}{\lambda}+\frac{3\lambda^2}{a^2\Phi_0^2}\right]\frac{H}{8\pi}, \tag{10}$$

where $\Phi_0\equiv\Phi/\Phi_\infty$ and $\Phi_\infty^2\equiv m_ec^2/4\pi e^2\lambda^2$.

After a somewhat lengthy computation (see Appendix), the reduced superheating field, $h_{sh}\equiv H_{sh}/H_c$, is given by

TABLE I. Critical properties of Sn and Al.

Material	T_c(K)	$H_c(0)$(G)	ξ_0 (μm)	λ (nm)
Sn	3.7	309	0.23	34
Al	1.2	105	1.6	16

$$h_{sh}=\left(1+\frac{4}{\sqrt[4]{15}}\gamma_G\right)h^0_{sh}, \tag{11}$$

where $\gamma_G\equiv 3\gamma/2\sqrt{|\alpha|\beta}$ is a dimensionless parameter, and h^0_{sh} is the "unperturbed" ($\gamma=0$) superheating field.

Generally, such measurements have been obtained with colloids, and the size distributions of the microspheres[17,18,20–22] provide a statistical error which renders a direct fit of γ_G from $h_{sh}(t)$ data unfeasible. Although the measurement reported in Ref. 19 was conducted using single microspheres, non-local and impurity effects lead to a large theoretical uncertainty in

$$h^0_{sh}(t)=\frac{1}{\sqrt{\kappa\sqrt{2}}}h^0_c(0)(1-t^2), \tag{12}$$

where $h^0_c(0)=H_c(0)/H_c(t)$, which is itself an approximation valid only close to T_c.[24]

Since the supercooling field implies the evaluation of a second derivative at the origin, it can be easily seen that the presence of a cubic term has no effect:

$$\frac{d^2}{d\Phi^2}(\gamma_G\Phi^3)=6\gamma_G\Phi\rightarrow 0. \tag{13}$$

The effect of γ (through γ_G) on the superheating field must be small, otherwise it would have been already detected; therefore, we must have $\gamma_G\ll 1$, which is not valid for relative temperatures in the range

$$1-\frac{9}{4}\frac{\delta^2}{\alpha'\beta}<t<1. \tag{14}$$

For this interval to be vanishingly small,

$$\frac{9}{4}\frac{\delta^2}{\alpha'\beta}\ll 1. \tag{15}$$

This is a weaker bound than the one of Eq. (9).

For Al and Sn with maximum critical fields of order 10^2 G, the shift between h_{sh} and h^0_{sh} is less than 10^{-2} G, well below the sensitivity of measurements.[17–22] For this reason, we simply drop the bound of Eq. (15) and consider only the one of Eq. (9). Conversely, a breakdown of the perturbation expansion of the superheating reduced field would imply a superheating temperature shifted to $t_{sh}=1.6$. Also notice that no "spikes" should be seen in the H-T superheating curve for values of t in the "exclusion" interval as the field values are quite small there.

Table II provides a comparison of the bounds on δ with the prediction of Ref. 13. The analogy between cosmology and condensed matter prompts for comparison with the TFT cubic term also. To do this, we compute the associated slope of $\gamma(t)$ from $\gamma(T)=(\sqrt{2}/4\pi)e^3T$, obtaining $\delta=(\sqrt{2}/4\pi)e^3T_c$. Obviously, although the prediction is material independent, its formulation in terms of a reduced temperature is not.

TABLE II. Derived quantities and bounds for δ.

Material		Sn		Al
α' (J)		1.15×10^{-25}		2.38×10^{-27}
β (J m^3)		4.72×10^{-54}		2.16×10^{-56}
α' (eV^2)		3.61×10^{-1}		7.45×10^{-3}
β		9.45×10^{-4}		4.32×10^{-6}
Bound	t_{sh} shift	δ (eV)	t_{sh} shift	δ (eV)
h_{sh}	0.25	1.23×10^{-2}	0.25	1.20×10^{-4}
ΔT_{exp}	10^{-3}	7.78×10^{-4}	10^{-3}	7.57×10^{-6}
Ref. 13	5.19×10^{-9}	1.77×10^{-6}	2.32×10^{-6}	3.64×10^{-7}
TFT	2.92×10^{-9}	1.33×10^{-6}	3.24×10^{-6}	4.31×10^{-7}

Note that the dimensionality of γ here is changed with respect to the free-energy potential of Eq. (5), through a convenient m_e factor—this is because the dimension of the scalar field in GL theory is $[\Phi^2]=L^{-3}$, its square representing a density, while in field theory $[\Phi]=L^{-1}$. The electron mass determines the conversion as it is absent from the kinetic term of the Lagrangean density of field theory, $\partial_\mu\Phi\partial^\mu\Phi$, while present in the corresponding condensed-matter free energy term, $(\hbar^2/2m_e)\nabla^2\Phi$ (or, equivalently, in the coherence length: $\xi^2_{FT}=1/\alpha'$ versus $\xi^2_{cm}=\hbar^2/2m_e\alpha'$).

Table II also includes the quantities α' and β, both in SI and natural units. As explained above, conversion is not direct, but achieved through the multiplicative factor m_e.

The results in Table II include the cubic term predicted by Ref. 13. In the absence of an applied magnetic field, each momentum fluctuation of the gauge field $\vec{A}$ has an expectation value given by the equipartition theorem. When suitably integrated over the momentum space (with a cutoff Λ of the order of ξ_0^{-1}),

$$\langle A^2\rangle_\Phi=4\frac{\mu_0}{\pi}\Lambda T_c-8\mu_0\frac{e}{\hbar c}\sqrt{\pi\mu_0}T_c|\Phi|. \tag{16}$$

Since A^2 couples to Φ^2 in Eq. (1), this translates into an unimportant correction to the scalar field mass, plus a negative cubic term, given by $-8\mu_0(e/\hbar c)\sqrt{\pi\mu_0}T_c|\Phi|^3$. This term implies a shift in the superheating temperature (at zero field), of

$$\Delta_T=7.25\times10^{-12}T_c^3H_c(0)^2\xi_0^6, \tag{17}$$

with $H_c(0)$ in G and ξ_0 in μm.

This shift lies beyond experimental accessibility, since it requires a temperature sensitivity of 10^{-6} K (for Al; 10^{-9} K for Sn). However, such an experiment *might* be performed with Al, using state-of-the-art relative temperature measurement techniques.

Surprisingly, for both materials the slopes of γ predicted by TFT and Ref. 13 have similar magnitudes, $\sim10^{-7}$ eV. This is an indication of the analogous underlying mechanisms behind them: the thermal averaging of the gauge field in condensed matter can be thought of as equivalent to finite-temperature vacuum polarization in high-energy physics (expressed by the renormalization of one-loop Feynmann diagrams).

IV. TOPOLOGICAL DEFECT FORMATION

Let us now discuss some possible implications of the inclusion of the cubic term in the mean-field potential. Since temperature sensitivity measurements constrain $\gamma_G < 10^{-2}$, we always assume $\gamma^2 \ll \alpha\beta$.

The K-Z mechanism predicts a density of topological defects (vortices), $n \simeq \xi_0^{-2}(\tau_0/\tau_q)^{\nu}$, where $\tau_0 = \pi\hbar/16k_BT_c$ is the characteristic time scale, given by the Gorkov equation, τ_q is the quench time, and ν is a critical exponent. Moreover, it rests upon the assumption that there is a single topological defect per ξ_0^2 area and, therefore, one must look for changes induced in this quantity. However, since the characteristic scales of the problem are obtained via linearization of the GL equations, close to T_c and when the order parameter is small, we see *no* changes in this prediction.

On the other hand, for a thin slab of of width L_z, the HR mechanism predicts a defect density of the order $n \simeq (e/2\pi)T^{1/2}L_z^{-1/2}\hat{\xi}^{-1}$, where $\hat{\xi} \sim 2\pi/\hat{k}$ is the domain size immediately after the transition. This quantity is related to the highest wave number $\hat{k}$ to fall out of equilibrium and is obtained from the adiabaticity relation

$$\left|\frac{d\omega(k)}{dt}\right| = \omega^2(k), \tag{18}$$

for a given dispersion relation $\omega(k)$. In the underdamped case, $\omega(k) = \sqrt{k^2+m_\gamma^2}$, with a photon mass given by $m_\gamma^2 = 2e^2|\Phi|^2 = -2e^2\alpha/\beta$. Thus we obtain $\hat{k} \sim \sqrt[3]{\alpha' e^2/\beta\tau_q}$, and hence $n \propto \tau_q^{-1/3}$. Here, the introduction of a cubic term in the potential will change the photon mass, as the true vacuum shifts to

$$\Phi = \frac{-3\gamma + \sqrt{-16\alpha\beta + 9\gamma^2}}{4\beta}. \tag{19}$$

However, since $\gamma_G \equiv 3\gamma/2\sqrt{|\alpha|\beta} \ll 1$, the effect of the cubic term is too small to significantly change the HR result.

Another effect related to metastability concerns the nonvanishing probability of the order parameter to quantum tunnel from the the symmetric (false) vacuum towards the nonsymmetric vacuum. Following Refs. 25 and 26, the rate of transition per unit volume and time to the true vacua is given, in the thin wall approximation, by

$$\frac{\Gamma}{V\Delta t} = T^4\left(\frac{S_3}{2\pi T}\right)^{3/2} e^{-S_3/T}, \tag{20}$$

where

$$S_3(T) = \frac{2\pi}{81}\frac{1}{\beta^7\sqrt{\beta}}\frac{\gamma^9(T)}{\epsilon^2(T)} \tag{21}$$

is the Euclidean action, and $\epsilon(T)$ is the "depth" of the true vacuum.

The thin wall approximation is valid whenever the barrier's height is much greater than ϵ. This is true when γ is comparable to the other parameters, namely, when $\gamma_G \sim 1$. For eventually smaller values of γ, like those predicted by TFT and Ref. 13, the approximation fails. In fact, we have shown that the current temperature sensitivity of 10^{-3} K only allows values of γ_G smaller than 10^{-2}. Therefore, we must conclude that the barrier's height is not comparable to the true vacuum's depth, and the field should always tunnel through it (i.e., with a probability close to unity). Because of this, there is no concern that defects may not have time to nucleate within the resolution time of the measuring device, as would happen if the potential barrier were high and diminished too slowly.

V. CONCLUSIONS

In this work we have examined a possible description of a type-I superconducting phase transition by introducing a cubic term in the GL mean-field potential, inspired by a gauge field thermal averaging,[13] and also by analogy with TFT.

Our analysis of the bounds derived from the superheating field and temperature constraints clearly shows that the contribution of any cubic term is small compared to other parameters in the GL potential. Thus the following conclusions can be drawn: First, the superheating temperature shift induced by a cubic term, derived either from Ref. 13 or from TFT, increases with decreasing G-L parameter κ [$\Delta t_{sh}(\mathrm{Sn}) \sim 10^{-9}$; $\Delta t_{sh}(\mathrm{Al}) \sim 10^{-6}$]. This suggests that future experiments to search for a TFT cubic term should be conducted with extreme ($\kappa \ll 1$) type-I materials, for example α-tungsten, with $T_c = 15.4 \pm 0.5$ mK, $H_c = 1.15 \pm 0.03$ G. Similarly, the shift in the supercritical field might be reinvestigated using a dc SQUID, which currently possess a sensitivity of $10^{-5}\phi_0/\sqrt{\mathrm{Hz}}$, or 10^{-6} G over a 10 μm grain diameter.

Furthermore, the impact of any cubic term on the defect density predictions of KZ (Refs. 1 and 5 or HR (Ref. 11) is negligible, with no suppression or slowing down of defect production because the potential barrier due to γ is not sufficient to prevent nucleation.

These considerations suggest that, all else being equal, experiments to detect topological defect formation in type-I superconductors would observe a reduction of the predicted HR defect densities by 10–100 depending on choice of material. Recent calculations.[27] however, suggest that the defect structure formed in type-I materials survives significantly longer than in type-II. Given that the type-I estimate is of order 10^{-4} s, it seems possible that the disadvantage in ξ might be compensated by simple measurability.

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APPENDIX

Including the presence of a magnetic field with a Φ dependent magnetic moment, the conditions for the reduced superheating field $h_{sh} \equiv H_{sh}/H_c$ are

(i) minimum

$$h_{sh}^2 = \frac{8}{9}\left[\frac{\Phi_0^2\ (1+\gamma_G\Phi_0-\Phi_0^2)\sinh^2(x_0)}{1+\dfrac{\sinh(2x_0)}{2x_0}-\dfrac{2\sinh^2(x_0)}{x_0^2}}\right],$$

(ii) inflexion point

$$h_{sh}^2 = \frac{4}{9}\left[\frac{\Phi_0^2\ (1+2\gamma_G\Phi_0-3\Phi_0^2)x_0^2\sinh^2(x_0)}{3\sinh^2(x_0)-x_0^3\coth(x_0)-x_0^2-\dfrac{1}{2}x_0\sinh(2x_0)}\right],$$

where $x_0 \equiv \Phi_0 a/\lambda$.

Since the diameter a is much larger than the penetration depth λ, we can take the limit $x_0 \to \infty$. The above conditions become

$$h_{sh}^2 = \frac{8}{9}\Phi_0^3\ (1+\gamma_G\Phi_0-\Phi_0^2)\frac{a}{\lambda}$$

and

$$h_{sh}^2 = \frac{4}{9}\Phi_0^3(1+2\gamma_G\Phi_0-3\Phi_0^2)\frac{a}{\lambda}.$$

Solving for Φ_0 we get

$$2(1+\gamma_G\Phi_0-\Phi_0^2) = -(1+2\gamma_G\Phi_0-3\Phi_0^2),$$

which implies that

$$\Phi_0 = \sqrt{\frac{3}{5}+\frac{4}{25}\gamma_G^2}+\frac{2}{5}\gamma_G \simeq \sqrt{\frac{3}{5}}+\frac{2}{5}\gamma_G.$$

Substituting in the first expression, we obtain

$$h_{sh}^2 = \frac{8}{9}\left(\sqrt{\frac{3}{5}}+\frac{2}{5}\gamma_G\right)^3\left[1+\gamma_G\left(\sqrt{\frac{3}{5}}+\frac{2}{5}\gamma_G\right) -\left(\sqrt{\frac{3}{5}}+\frac{2}{5}\gamma_G\right)^2\right]\frac{a}{\lambda},$$

which, to first-order in γ_G, becomes

$$h_{sh} = \frac{4}{5\sqrt[4]{15}}\left(1+\frac{4}{\sqrt[4]{15}}\gamma_G\right)\sqrt[4]{\frac{a}{\lambda}},$$

giving the result Eq. (11), with

$$h_{sh}^0 = \frac{4}{5\sqrt[4]{15}}\sqrt[4]{\frac{a}{\lambda}}.$$

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