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Quantum tunneling of flux lines in a high-T_c superconductor

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The quantum flux creep model of Ivlev *et al.* has been extended to include the dissipative and inertial mass contributions to the bounce action. The crossover temperature, the activation energy, and the quantum tunneling rate are evaluated numerically. This has been carried out by a perturbation expansion retaining only the leading-order term deemed to be valid in the vicinity of and below the crossover temperature. It is found that the crossover temperature and the activation energy are relatively unchanged but the quantum tunneling rate is drastically reduced when compared with the estimates of Ivlev *et al.* Furthermore, the quantum tunneling rate is temperature dependent in the neighborhood of the crossover temperature.

I. INTRODUCTION

The phenomena of quantum tunneling of flux lines in the mixed state of superconductors at very low temperatures has attracted a great deal of attention recently. The first experimental observations were made by Mitin[1] in his study of the relaxation of the flux-line gradients in the PbMoS system below 1 K and the data were interpreted as signaling the occurrence of quantum tunneling of vortices between pinning centers. Low-temperature magnetic relaxation studies of the vortex state have been carried out in many superconducting systems, including the high-T_c superconductors,[2] LaSrCuO, high-quality single crystals[3] of YBCO, organic superconductors[4] and heavy-fermion systems.[5] Typical results from these experiments are summarized in Ref. 6.

According to these experiments, the magnetic relaxation rate $S=(\partial \ln M)/(\partial \ln t)$ exhibits two distinct types of behavior as a function of temperature. Above a characteristic temperature T_0 the decay rate is of the Arrhenius type. This behavior is interpreted as arising from the classical "jump over the hump" motion of flux lines and is characterized by an activation energy. Below T_0 the crossover temperature, the decay rate is essentially independent of temperature, and is interpreted as arising from the quantum tunneling of vortices.

There has also been a considerable amount of theoretical activity in this area. Blatter, Geshkenbein, and Vinokur[7] have calculated the tunneling rates for single vortices and vortex bundles in bulk isotropic superconductors within the framework of weak collective pinning theory. It was later extended to the case of anisotropy and layered superconductors by Blatter and Geshkenbein.[8] The theoretical analysis and the numerical estimates given in Refs. 7 and 8 are based on dimensional analysis. Therefore, only the limiting values as $T\rightarrow 0$ of the theoretical tunneling rate are compared with the experimental data Ref. 6. No detailed calculations of the tunneling rate at nonzero temperatures could be made.

The quantum creep of vortices in layered superconductors was studied by Ivlev, Ovchinnikov, and Thompson,[9] who determined the crossover temperature and the activation energy as a function of the transport current, the normal resistivity, and the upper critical field. No detailed numerical results for quantum tunneling were presented by them either. Ma *et al.*[10] do claim a more exact treatment of the activation energy, quantum tunneling, and flux-creep rates on the basis of a quartic pinning potential model. However, they have not included the Lorentz potential in their model. Therefore, their calculation does not address the issue of the current and field dependence of the various physical quantities. It has been emphasized by Coffey and Clem[10(a)] that inertial mass may be important in the dynamics of vortices in high-T_c superconductors. While Blatter *et al.*[7,8] and Ma *et al.*[10] do include the inertial mass term, Ivlev, and Ovchinnikov, and Thompson[9] do not include it in their treatment. Moreover, all the works referred to (Refs. 7–10) consider the quantum decay rate only in the limit of $T\rightarrow 0$. While this approach may be adequate for systems with crossover temperature, T_0, in the mK range, it is not for systems with relatively large values of T_0. For example, Ivlev, Ovchinnikov, and Thompson,[9] estimate T_0 to be $\sim$4 K but Ma *et al.*[10] obtain $T_0 \sim 40$ K. Although the primary mode of decay below T_0 is due to tunneling, the decay rate is not independent of temperature. Furthermore, when T_0 is sufficiently high, finite temperature effects due to fluctuations may also become important. Such effects have not been addressed. Moreover, the method of Ref. 9 which includes only the first harmonic in the bounce is not adequate for treating the temperature dependence of the quantum tunneling rate, which is determined by the spectral characteristics of the coupling to the environment.

It is the purpose of this paper to carryout the analysis of the decay rate near and below the crossover temperature T_0 and to examine its dependence on temperature, current, and field variables. The general method based on functional integrals for treating the influence of dissipative processes on the decay of quantum systems from metastable states via quantum tunneling was described by Caldeira and Leggett.[11] In this paper we will consider the case of straight flux lines in a layered high-T_c superconductor. An outline of the theoretical framework will

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be given in Sec. II. The present model represents an improvement of the model of Ref. 9 by including (i) the contributions of the inertial mass term and the viscous damping term in calculating the action, (ii) all the higher-order harmonics in the expression for the bounce, and (iii) by extending beyond the harmonic approximation in the pinning potential. The resulting equations are quite complex to obtain analytic solutions through the temperature range $0-T_0$. A perturbation analysis near the crossover temperature yields expressions for the crossover temperature, activation energy, and the tunneling rate and its dependence on current and temperature. It will be seen that the crossover temperature and the activation energy are not greatly influenced by these modifications, the values are essentially the same as those given by Ivlev, Ovchinnikov, and Thompson.[9] But the quantum tunneling rate due to the bounce contribution[6] is substantially reduced and it is temperature dependent in the neighborhood of T_0, $T<T_0$. These limitations and possible extensions will be discussed in the final section.

II. THEORETICAL FRAMEWORK

In this paper we will follow mainly the notation and the theoretical frame work of Ivlev, Ovchinnikov, and Thompson[9] but include the effect of the inertial mass term and the damping term. We will also include the higher-order harmonics in the bounce and we go beyond the harmonic approximation in the pinning potential. For the sake of clarity we explicitly retain $\hbar$ and use $\beta=1/k_B T$ in the equations.

A. The effective action

Consider a straight flux line of mass M per unit length, moving in a potential $V(u)$ and undergoing elastic deformation characterized by the displacement u. The dissipative effects are included by coupling to a heat-bath environment represented by a set of harmonic oscillators.[11] The magnetic field is applied in the y direction. The vortex displacement $u(y,t)$ has only a z component. For the semiclassical treatment of the motion of the flux line, we consider the effective Euclidean action ($\tau=-it$) given by

$$A=\int_{-\infty}^{\infty}dy\int_0^{\hbar\beta}d\tau\left[E_{\rm el}\left[\frac{\partial u}{\partial y}\right]+V(u)+E_D\left[\frac{\partial u}{\partial \tau}\right]+\frac{1}{2}M\left[\frac{\partial u}{\partial \tau}\right]^2\right] . \tag{1}$$

Here $E_{\rm el}(\partial u/\partial y)=(\varepsilon/2)(\partial u/\partial y)^2$ is the elastic term given in terms of the line tension ε of the flux-line lattice,

$$\varepsilon=\frac{\phi_0^2}{8\pi^2}\frac{\lambda_c}{\lambda_{ab}^3}\ln\left[\frac{\lambda_c}{\xi_{ab}}\right] .$$

$V(u)$ is the pinning potential consisting of a periodic part and the Lorentz potential:

$$V(u)=-\frac{\phi_0 j_c d}{2\pi c}\cos\left[\frac{2\pi u}{d}\right]-\frac{\phi_0 j}{c}u . \tag{2}$$

E_D, the so-called Caldeira-Leggett action, describes the phenomenological coupling to the heat-bath environment and can be written as

$$E_D\left[\frac{\partial u}{\partial \tau}\right]=-\frac{\eta}{2\pi}\frac{\partial u}{\partial \tau}\int_0^{\hbar\beta}d\tau_1\frac{\partial u}{\partial \tau_1}\ln\left|\sin\frac{\pi}{\hbar\beta}(\tau-\tau_1)\right| . \tag{3}$$

B. The crossover temperature

We consider the limit of a large current which is only slightly less than the critical depinning current, i.e., $j_c-j\ll j_c$ then the potential in Eq. (2) can be expanded around the inflection point and omitting the trivial shift constants it can be written as

$$V(u)=\frac{\phi_0 j_c d}{2\pi c}\left[\left[\frac{2\pi u}{d}\right]^2\left[\frac{j_c-j}{2j_c}\right]^{1/2}-1/6\left[\frac{2\pi u}{d}\right]^3\right] . \tag{4}$$

Obviously the potential is zero at $u=0$ and at $u_m=(3d/\pi)[(j_c-j)/2j_c]^{1/2}$. The distance u_m will be used for scaling purposes.

The quantum-mechanical tunneling probability is determined by the action in Eq. (1). In the semiclassical approximation it is given by $W\sim\exp(-A_0/\hbar)$, where A_0 is the value of the action on a classical trajectory for which the equation of motion is

$$-M\frac{\partial^2 u}{\partial \tau^2}-\varepsilon\frac{\partial^2 u}{\partial y^2}+V'(u)+\frac{\eta}{\hbar\beta}\int_0^{\hbar\beta}d\tau_1\frac{\partial u}{\partial \tau_1}\cot\frac{\pi(\tau-\tau_1)}{\hbar\beta}=0 . \tag{5}$$

According to Ivlev, Ovchinnikov, and Thompson,[9] at high temperatures the classical solution $u(y,\tau)$ does not depend on τ. The action A_0 is determined by the static function $u_0(y)$ obeying the equation

$$-\varepsilon\frac{\partial^2 u_0}{\partial y^2}+V'(u_0)=0 , \tag{6}$$

and the activation energy U_0 is given by $A_0=U_0\hbar\beta$. There is a crossover temperature, T_0, above which there is a thermally activated regime. Below T_0, a new semiclassical trajectory called the bounce develops. This is periodic in imaginary time. The crossover temperature is determined from the bounce in the following way. The displacement just below T_0 can be written as a Fourier expansion around the static solution u_0 of Eq. (6):

$$u(y,\tau)=u_0(y)+\sum_{n=1}^{\infty}\psi_n(y)\cos(\omega_n\tau) . \tag{7}$$

Here $\omega_n=2\pi n/\hbar\beta$ are the Matsubara frequencies. Substituting from Eq. (7) into the equation of motion (5), one obtains the following equation for $\psi_n(y)$:

$$-\varepsilon\frac{\partial^2\psi_n}{\partial y^2}+V''(u_0)\psi_n=-(\eta\omega_n+M\omega_n^2)\psi_n . \tag{8}$$

If one introduces new variables[9]

$$v=\frac{u}{u_m},\quad \phi_n=\frac{\psi_n}{u_m},$$

and

$$\zeta=\frac{y}{y_0},$$

where

$$y_0^2=\frac{\varepsilon dc}{\pi\phi_0 j_c}\left[\frac{2j_c}{j_c-j}\right]^{1/2},\tag{9}$$

Eq. (6) becomes

$$-\frac{1}{2}\frac{\partial^2 v_0}{\partial\zeta^2}+2v_0-3v_0^2=0\tag{10}$$

with the solution

$$v_0=\operatorname{sech}^2\zeta\ .\tag{11}$$

Equation (8) turns into

$$-\frac{1}{2}\frac{\partial^2\phi_n}{\partial\zeta^2}+2(1-3\operatorname{sech}^2\zeta)\phi_n=E_n\phi_n\ ,\tag{12}$$

where

$$E_n=-\frac{dc}{2\pi\phi_0 j_c}\left[\frac{2j_c}{j_c-j}\right]^{1/2}(\eta\omega_n+M\omega_n^2)\ .\tag{13}$$

Now Eq. (12) has three discrete eigenvalues $-\frac{5}{2}$, 0, and $\frac{3}{2}$, with the associated (unnormalized) eigenfunctions, $\operatorname{sech}^3\zeta$, $\operatorname{sech}^2\zeta\tanh\zeta$, and $(5\tanh^2\zeta-1)\operatorname{sech}\zeta$, respectively. The crossover temperature is determined by the negative eigenvalue of Eq. (12). Therefore the crossover temperature is determined by using Eq. (13) after substituting for $\omega_1=2\pi/\hbar\beta_0$:

$$-\frac{5}{2}=-\frac{dc}{\phi_0 j_c}\left[\frac{2j_c}{j_c-j}\right]^{1/2}\left[\frac{\eta}{\hbar\beta_0}+\frac{2\pi M}{(\hbar\beta_0)^2}\right].\tag{14}$$

One can solve for the positive root for T_0:

$$k_B T_0=\frac{1}{\beta_0}=\frac{-\hbar\eta}{4\pi M}+\frac{\hbar\eta}{4\pi M}\left[1+\frac{10\phi_0 j_c}{dc}\frac{2\pi M}{\eta^2}\left[\frac{j_c-j}{2j_c}\right]^{1/2}\right]^{1/2}.\tag{15}$$

Equation (15) gives the crossover temperature as a function of the parameters M, η, and j.

C. The activation energy

In the thermally activated region, i.e., above T_0, the activation energy can be obtained by evaluating the action defined in Eq. (1) for the static solution as $A_0=U_0\hbar\beta$. When the static solution of Eq. (6) is used, the contribution to the action A_0 arises only from the elastic and pinning terms in Eq. (1) and there is no contribution from the dynamic terms involving the inertial mass and the viscous damping. The action is then given by

$$A_0=\frac{48}{5\pi}d\left[\frac{\varepsilon d\phi_0 j_c}{\pi c}\right]^{1/2}\left[\frac{j_c-j}{2j_c}\right]^{5/4}\hbar\beta\ .\tag{16}$$

The activation energy U_0 is given by the coefficient of $\hbar\beta$ in Eq. (16). This result is identical to that obtained in Ref. 9.

III. QUANTUM TUNNELING

Below T_0, the predominant mechanism by which the metastable state decays is quantum-mechanical tunneling. The rate at which this takes place is determined by the action for the bounce trajectory. Ivlev, Ovchinnikov, and Thompson[9] have determined the Euclidean action applicable to this case as

$$A_0^q=\frac{48}{5\pi}d\left[\frac{\varepsilon d\phi_0 j_c}{\pi c}\right]^{1/2}\left[\frac{j_c-j}{2j_c}\right]^{5/4}\hbar\beta_0\tag{17}$$

and hence according to them, the quantum tunneling rate $\sim\exp[-A_0^q/\hbar]$ is essentially independent of temperature. However, it should be noted that this result holds good for the saddle-point trajectory corresponding to the static solution; it includes the contributions from the elastic and pinning terms only and is true only in the limit of $T\to T_0$. For temperatures in the range $T<T_0$ the bounce trajectory will be different from the saddle-point trajectory of the static solution and the contributions of the inertial mass term and the viscous damping term will also have to be included in the bounce action. In addition one would have to go beyond the harmonic approximation in the pinning potential as employed in Eq. (12):

$$v'(u)=v'(u_0)+v''(u_0)(u-u_0)+\tfrac{1}{2}v'''(u_0)(u-u_0)^2\ .\tag{18}$$

The new equation of motion in terms of dimensionless quantities becomes

$$-\frac{1}{2}\frac{\partial^2\phi_n}{\partial\zeta^2}+2(1-3\operatorname{sech}^2\zeta)\phi_n-E_n\phi_n=3\sum_{m-1}^{\infty}\phi_{n+m}\phi_m+\frac{3}{2}\sum_{m=1}^{n}\phi_{n-m}\phi_m\ ,\tag{19}$$

where E_n is still given by Eq. (13).

While an analytical solution of Eq. (19) is not possible, numerical solution by successive iteration[12,13] can be obtained over a wide range of temperature and mass and damping parameter values. Moreover, very near the crossover temperature, a perturbation analysis in terms of a parameter $\lambda=[(T_0-T)/T_0]^{1/2}$ is possible in view of the temperature dependence of the terms in E_n, Eq. (13). It can be shown that ϕ_n are of the order λ^n and that the leading-order contribution arises from $\phi_1\sim\operatorname{sech}^3\zeta$. Substituting for u, with the leading-order approximation from Eq. (7), we have evaluated the action from Eq. (1). The revised contribution from the elastic and pinning terms to the action of the bounce is given by

$$A_p^q=\frac{438}{35}d\left[\frac{\varepsilon d\phi_0 j_c}{\pi c}\right]^{1/2}\left[\frac{j_c-j}{2j_c}\right]^{5/4}\hbar\beta\ . \tag{20}$$

The contribution from the inertial mass term is

$$A_m^q=\frac{48Md^2}{5}\left[\frac{\varepsilon dc}{\phi_0 j_c\pi}\right]^{1/2}\left[\frac{j_c-j}{2j_c}\right]^{3/4}\frac{1}{\hbar\beta} \tag{21}$$

and the contribution from the viscous damping term is

$$A_d^q=\frac{24\eta d^2}{5\pi}\left[\frac{\varepsilon dc}{\phi_0 j_c\pi}\right]^{1/2}\left[\frac{j_c-j}{2j_c}\right]^{3/4}\ . \tag{22}$$

The total action for the bounce trajectory is given by $A^q=A_p^q+A_m^q+A_d^q$. Finally, the quantum tunneling rate is given by the standard WKB formula:

$$\Gamma\sim(A^q/2\pi\hbar)^{1/2}\exp(-A^q/\hbar)\ . \tag{23}$$

IV. RESULTS AND DISCUSSION

We have derived expressions for the bounce contribution to the action arising from the elastic and pinning terms, Eq. (20), from the inertial mass term Eq. (21), and from the damping term Eq. (22) by retaining only the leading-order term in the bounce. The last two are new results and Eq. (20) differs by a numerical factor from the result of Ivlev, Ovchinnikov, and Thompson.[9] We have also calculated numerically the crossover temperature T_0 from Eq. (15), the activation energy from Eq. (16), and the exponent $A^q/\hbar$ as a function of temperature. The values of the input parameters are listed in Table I and are the same as the ones used by Ivlev, Ovchinnikov, and Thompson[9] in their estimates.

Figure 1 shows the variations of T_0 with the fractional current difference (FCD), $(j_c-j)/j_c$. Also shown are the values calculated by using the expression given by Ivlev, and Ovchinnikov, and Thompson.[9] As can be seen the two curves are practically identical; T_0 is about 4 K for a value of 0.0035 for FCD and increases to 21 K for a value of about 0.1 for FCD.

Figure 2 shows the variation of the activation energy as a function of FCD. This result is also identical with that obtained by Ivlev, Ovchinnikov, and Thompson.[9] The activation energy increases from 0.2 meV at 0.0035 for FCD to about 13 meV for a value of 0.1 for FCD.

Figure 3 shows the variation of action for the bounce trajectory with temperature (solid line). Also shown is the contribution from only the elastic and pinning potential terms (thin dashed line) for comparison with the Ivlev expression (thick dashed line). The Ivlev result ignores the mass and inertial terms. It also ignores the contribution to the elastic and pinning terms from the dynamic term of the bounce in Eq. (7). Finally, Fig. 4 shows the variation of the function $\exp(-A^q/\hbar)$ with temperature. Again the thick dashed line corresponds to the Ivlev approximation.

It is obvious that the crossover temperature and the activation energy are not significantly influenced by the mass and the damping terms. Above the crossover temperature the decay rate is mainly determined by the saddle-point trajectory of the static solution. There are corrections to the activation energy above T_0, but these are beyond the scope of the present work.

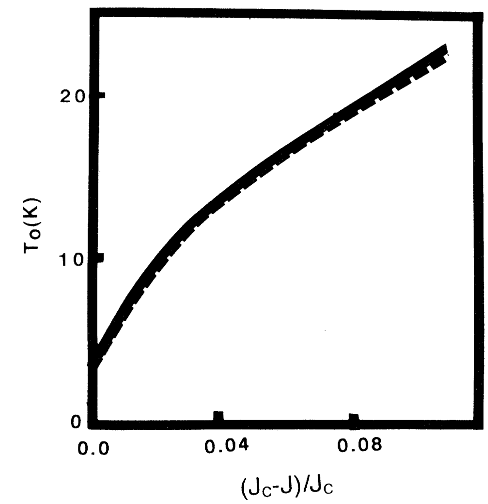

FIG. 1. Variation of the crossover temperature with current: (a) solid line, present work; (b) dashed line, Ivlev, Ovchinnikov, and Thompson (Ref. 9).

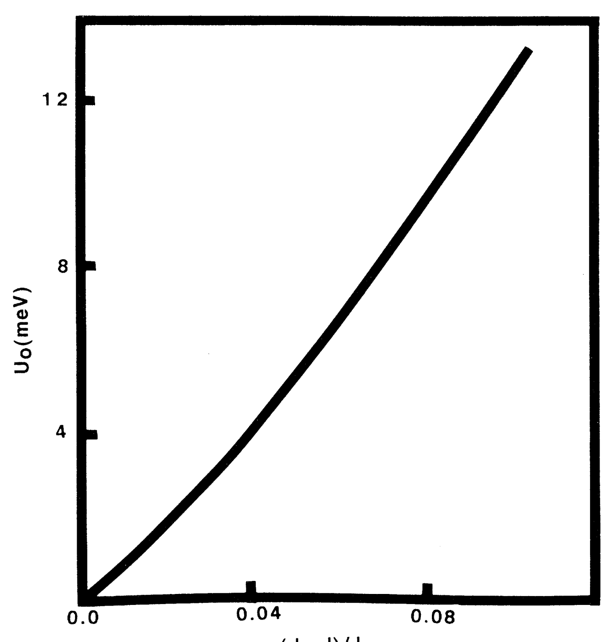

FIG. 2. Variation of the activation energy with current.

TABLE I. Input parameters.

d: 1 nm	η: 4.12×10^{-7} ns/m^4
Hc_2: 200 T	$\lambda_c/\lambda_{ab}=\xi_{ab}/\xi_c=5$
j_c: 1.0×10^{12} A/m^2	$\xi_{ab}=100$ nm
M[a]: 2.73×10^{-22} Kg/m	$\xi_c=1.0$ nm

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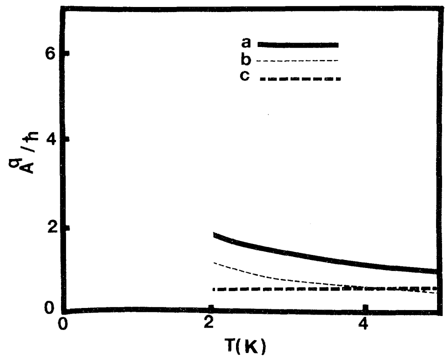

FIG. 3. Variation of the action for the bounce trajectory with temperature: (a) present work; (b) Ivlev modified by bounce action; (c) Ivlev, Ovchinnikov, and Thompson (Ref. 9).

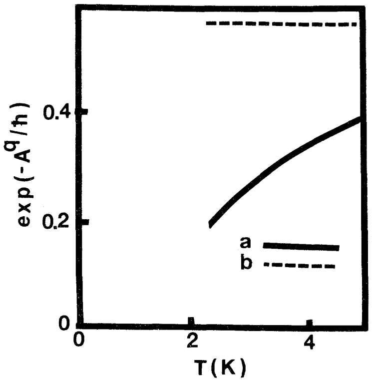

FIG. 4. Variation of $\exp(-A^q/\hbar)$ with temperature: (a) present work; (b) Ivlev, Ovchinnikov and Thompson (Ref. 9).

An important effect of including the contribution from the bounce trajectory is that the action is rendered temperature dependent and hence the quantum tunneling rate becomes temperature dependent below the crossover temperature. The effect of including the damping and inertial mass terms is to further reduce the quantum tunneling rate. One can give a hand-waving explanation of the explicit temperature variation of the different terms. The action integral is an integral over $d\tau$, the upper limit being the "bounce length" $\hbar\beta$. The integrand of the potential-energy term does not depend on τ, so it is simply multiplied by the factor $\hbar\beta$ as a result of the integration. The damping term, as it involves the first derivative with respect to τ, contains a factor $\omega_1 = 2\pi/\hbar\beta$; this factor cancels the $\hbar\beta$ from integration. The net result is that the damping term has no explicit dependence on temperature. The kinetic-energy term, on the other hand, involves the square of the first derivative and hence the factor, $\omega_1^2 = (2\pi/\hbar\beta)^2$. Integration yields a factor $\hbar\beta$, with the net result the factor $\hbar\beta$ in the denominator. These are the results obtained in Eqs. (20)–(22). We have verified that our equations do reproduce the well-known limiting values of activation energy in the limit $T \to T_0$ in all the special cases such as $M=0$.

The approach used here with a single harmonic for the bounce, however, is adequate for temperatures below T_0 and very near T_0. It is unsuitable for exploring the $T \to 0$ limit because the first-order perturbation expansion in Eq. (7) with a single harmonic is valid only near T_0. Therefore, one should not infer that Eq. (20) leads to a divergent behavior as $T \to 0$. An Arrhenius-like behavior has also been obtained for temperatures immediately below T_0 for a massive damped particle tunneling in a cubic potential. A correction to the Arrhenius behavior is obtained when the perturbation expansion is carried to the fourth order. Even higher-order expansions would be required to reach the $T=0$ limit. In any case, the quantum tunneling rate would be lower than the one given by Ivlev, Ovchinnikov, and Thompson. But, then the simplicity of the perturbation theory is lost. Quantum decay for the $T=0$ case has been studied analytically for several limiting cases and the dimensional estimates are well known.[14–16] However, for describing the finite-temperature effects over the entire temperature range, one would also have to include fluctuation modes in the bounce and finite order perturbation may not be possible. An alternative recourse is numerical analysis. These issues will be addressed in a future publication.

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