Effects of nonmagnetic scatterers on the magnetic coherence length in the mixed state of type-II superconductors

R. Laiho,* M. Safonchik,† and K. B. Traito
Wihuri Physical Laboratory, University of Turku, FIN-20014 Turku, Finland
(Received 18 March 2008; revised manuscript received 6 June 2008; published 29 August 2008)

Influence of impurities on magnetic coherence length ξ_h in mixed state of type-II superconductors is investigated in framework of quasiclassical theory. Nonmonotonic magnetic-field dependence of $\xi_h(B)$ is found. Increasing of the scattering rate results in decreasing of ξ_h making the $\xi_h(B)$ curve flat. Interconnection between ξ_h and the coherence length in the Meissner state is obtained. Characteristic relaxation time τ_0 is the only parameter needed for excellent fitting of the numerically calculated dependences of $\xi_h(\tau,B,T)$. The $\tau_0(B,T)$ dependence is calculated.

DOI: 10.1103/PhysRevB.78.064521 PACS number(s): 74.20.Rp, 74.20.Fg, 74.25.Op

Various characteristic lengths connected with space distribution of the order parameter and field distribution in the mixed state of type-II superconductors are discussed in Ref. 1. Recently it has been suggested that the magnetic properties of flux-line lattice (FLL) can be explained using only one fitting parameter—a field-dependent effective coherence length $\xi_h(B)$.[2] The dependence $\xi(B)$ is related to deviations in $M(\ln B)$ from linear behavior prescribed by London model. An empirical method to extract $\xi(B)$ from fitting of the magnetization data is proposed. It has been found from the comparison with experimental magnetization data in different materials that the perfect fitting can be obtained. Using this method influence of scattering on the shape of the $\xi(B)$ dependence has also been investigated in a series of $Lu(Ni_{1-x}Co_x)_2B_2C$ crystals in which the mean-free path is progressively reduced when increasing the Co content.[3] The results show flattening of the shape of the $\xi(B)$ dependence along with increasing relaxation time. Because this empirical method is successful, it is important to obtain microscopical justification for it. Effects of impurity scattering existing in real samples should be included in the model.

There are different theoretical approaches based on quasiclassical equations for describing the influence of impurities and temperature on magnetic-field dependence of the coherence length. The field dependence of $\xi(B)$ has been obtained analytically,[4] but the vector potential and the superconducting gap were not calculated self-consistently. The Helfand-Werthamer linearization technique for calculation of the vortex core size was applied. This technique is reasonable in the region of saturation of the pair potential $\Delta(r)$ characterized by length scale ξ_2 (see Table I of Ref. 1). However in this method the field distribution near the vortex center cannot be calculated accurately and the role of Kramer-Pesch effect[5] is not clarified. This theory predicts weakening of the field dependence of the core size ratio $\xi(B)/\xi(B_{c2})$ with increasing scattering.

Field dependence of the order-parameter coherence length ξ_Δ determined as $1/\xi_\Delta=[\partial|\Delta(r)|/\partial r]_{r=0}/|\Delta_{NN}|$ was calculated in Ref. 6. Here $|\Delta_{NN}|$ is the maximum value of the order parameter along the nearest-neighbor direction which is the direction of taking the derivative. In this theory $\xi_\Delta(B)/\xi_\Delta(B_{c2})$ increases with increasing scattering. However, the cutoff size extracted from magnetization data is not necessarily the same as the core size proportional to the slope of the order parameter at the vortex axis; approaching the core from outside to determine the cutoff, one may have a different result than when examining the core structure starting from its center.[2]

We have derived,[7] using self-consistent solution of quasiclassical Eilenberger equations for clean superconductors, an effective London model with the effective magnetic coherence length $\xi_h(B)$ as a parameter. In this approach the coherence length is obtained from the Ginzburg-Landau model[8,9] extended over the whole temperature range. A similar method was used for calculation of the Maki parameters $\kappa_1(T)$ and $\kappa_2(T)$ using the Eilenberger equations to distinguish the temperature dependences of the upper critical field H_{c2} and the initial slope of the magnetization M/H, respectively.[10] Our method describes the magnetic-field distribution both in the core region of rapidly growing $\Delta(r)$ and in the saturation area far from the core. Direct connection between $\xi_h(B)$ and magnetic-field distribution can be obtained from generalized London equation with the relation

$$h(\mathbf{r}) = \frac{\phi_0}{S}\sum_{\mathbf{G}} \frac{F(G)e^{i\mathbf{G}\mathbf{r}}}{1+\lambda^2 G^2}, \quad (1)$$

where $F(\mathbf{G})=uK_1(u)$, $K_1(u)$ is the modified Bessel function, $u=\sqrt{2}\xi_h G$, $\mathbf{G}$ is a reciprocal-lattice vector, and S is the area of the vortex lattice unit cell. This equation is valid in an intermediate magnetic-field range below the upper critical field. The aim of our paper is to calculate field corrections to usual London equation taking into account the impurities. The parameter ξ_h is responsible for the core effects in the mixed state and is different from the electromagnetic coherence length ξ_{el} in the Meissner state.[11] Experimentally, the $F(G)$ cutoff function can be investigated by small-angle neutron-scattering (SANS) measurements of FLL reflectivity and calculated from the integrated intensity of the Bragg peaks as the sample is rotated through the diffraction condition. Recently, the FLL form factor in $CeCoIn_5$ was found to be independent of the applied magnetic field, in striking contrast to the exponential decrease usually observed in superconductors.[12] This result is consistent with a strongly field-dependent coherence length, proportional to the vortex separation. It was proposed that the field-dependent coherence length is prominently observed in $CeCoIn_5$ due to the combination of large κ and the very high cleanliness of the sample.[12] Our previous calculation of field distribution in superclean superconductors also showed strong decrease in $\xi_h(B)$ in increasing moderate fields, followed by a minimum and growing again in higher fields.[7] It has been found that the symmetry of the paring state (s or d wave) is not crucial for the presence of the minimum in $\xi_h(B)$ dependence. The high-field regime in our model can be explained by the Abricosov solution of the Ginzburg-Landau theory which is not connected with any microscopical details of the model. This result can also be obtained from phenomenological consideration.[8] SANS measurements[12] were limited to moderate magnetic fields and did not reach the high-field limit which complicates comparison between these experiments and our theoretical approach. In the present paper we check the sensitivity of the $\xi_h(B)$ dependence to impurity scattering and find the range of relaxation times τ where the minimum in $\xi_h(B)$ exists.

We consider the mixed state of type-II superconductors at different levels of impurity scattering, looking for change in the shape of the $\xi_h(B)$ dependences for comparison with clean superconductors.[7] The magnetic-field penetration depth $\lambda(T)$ is assumed to be field independent and have the same value as in the Meissner state. However, $\lambda(T)$ is renormalized by impurity scattering[11,13] so that

$$\frac{\lambda^2}{\lambda_{L0}^2} = \left(2\pi T \sum_{n\geq 0} \frac{1}{\tilde{\Delta}_n (1+u_n^2)^{3/2}}\right)^{-1}, \tag{2}$$

where

$$\tilde{\Delta}_n = \Delta + \frac{1}{\tau\sqrt{u_n^2+1}}, \tag{3}$$

λ_{L0} is the London penetration depth at $T=0$ K, $u_n = w_n/\Delta$, the scattering time τ is given in units of $1/(2T_c)$, and $\omega_n = (2n+1)\pi T$.

We solve the quasiclassical self-consistent Eilenberger equations for the s-wave pairing symmetry. In what follows, the energy, the temperature, and the length are measured in units of T_c and the coherence length $\xi_0 = \xi_{BCS}\pi\Delta/T_c = v_F/T_c$. Hence $\xi_{BCS} = v_F/\pi\Delta$, where v_F is the Fermi velocity and Δ is temperature dependent uniform gap. The magnetic field $\mathbf{h}$ is given in units of $\phi_0/2\pi\xi_0^2$. In computations the ratio $\kappa = \lambda_{L0}/\xi_0 = 10$ is used. With the Riccati transformation of the Eilenberger equations[14,15] quasiclassical Green's functions f and g can be parametrized via functions a and b,

$$\bar{f} = \frac{2a}{1+ab}, \quad f^{\dagger} = \frac{2b}{1+ab}, \quad g = \frac{1-ab}{1+ab}, \tag{4}$$

satisfying the nonlinear Riccati equations,

$$\mathbf{u}\cdot\nabla a = -a[2(\omega_n + G) + i\mathbf{u}\cdot\mathbf{A}] + (\Delta + F) - a^2(\Delta^* + F^*), \tag{5}$$

$$\mathbf{u}\cdot\nabla b = b[2(\omega_n + G) + i\mathbf{u}\cdot\mathbf{A}] - (\Delta^* + F^*) + b^2(\Delta + F), \tag{6}$$

where the impurity potentials are determined as $F = \langle f\rangle/\tau$ and $G = \langle g\rangle/\tau$. One can expect the nonmonotonic magnetic-field dependencies of ξ_h obtained for clean superconductors[7] to remain in the limit of moderate scattering. Then the perturbation theory is valid and the Born approximation can be used for treating the impurity effects. This approach is considered in Eqs. (5) and (6). In the calculations we assume cylindrical Fermi surface and define $\mathbf{u}$ as unit vector of the Fermi velocity. The averages over isotropic Fermi surface can be reduced to averages over the polar angle φ, $\langle\cdots\rangle = (1/2\pi)\int\cdots d\varphi$.

To take into account the influence of screening the vector potential $\mathbf{A}(\mathbf{r})$ in Eqs. (5) and (6) is obtained from the equation

$$\nabla\times\nabla\times\mathbf{A} = \frac{4}{\kappa^2}\mathbf{J}, \tag{7}$$

where the supercurrent $\mathbf{J}(\mathbf{r})$ is given in terms of $g(\omega_n,\theta,\mathbf{r})$ by

$$\mathbf{J}(\mathbf{r}) = 2\pi T\sum_{\omega_n>0}\int_0^{2\pi}\frac{d\theta}{2\pi}\frac{\hat{\mathbf{k}}}{i}g(\omega_n,\theta,\mathbf{r}). \tag{8}$$

Here $\mathbf{A}$ and $\mathbf{J}$ are measured in units of $\phi_0/2\pi\xi_0$ and $2ev_F N_0 T_c$, respectively, and $\kappa = \lambda(T=0)/\xi_0$. The spatial variation of the internal field $h(\mathbf{r})$ is determined through

$$\nabla\times\mathbf{A} = \mathbf{h}(\mathbf{r}), \tag{9}$$

where $\mathbf{h}$ is measured in units of $\phi_0/2\pi\xi_0^2$.

The self-consistent condition for the pairing potential $\Delta(\mathbf{r})$ is given by

$$\Delta(\mathbf{r}) = VN_0 2\pi T\sum_{\omega_n>0}\int_0^{2\pi}\frac{d\theta}{2\pi}f(\omega_n,\theta,\mathbf{r}), \tag{10}$$

where V is the pairing interaction energy. The product VN_0 can be obtained from the expression

$$\frac{2}{VN_0} = \ln\frac{T}{T_c} + 2\pi T\sum_{0<\omega_n<\omega_c}\frac{1}{|\omega_n|}. \tag{11}$$

To obtain quasiclassical Green's function, the Riccati equations [Eqs. (5) and (6)] are solved by the fast Fourier transform (FFT) method. Unlike the square vortex lattice studied in Ref. 14, we consider a triangular vortex lattice, for which the wave vector mesh is transformed from square to hexagonal shape.[16]

Iterations of the coefficients a in the Fourier space are made using the relation

$$a(\mathbf{Q}) = \frac{[\Delta + F - (\Delta^* + F^*)a^2 - i\mathbf{u}\mathbf{A}]_{FT}}{i\mathbf{u}\mathbf{Q} + 2(\omega_n + G)}, \tag{12}$$

which makes it possible to solve the Riccati equations simultaneously for a full set of reciprocal vectors $\mathbf{Q}$ in the mesh, increasing greatly the calculation speed. The shifts to next iteration have to be damped[14] by a value depending on the

Matsubara frequency ω_n and the wave vector direction $\mathbf{u}$ forming the $\mathrm{damp}_{\mathrm{nu}}$ matrix. Every value in this matrix is optimized dynamically before the first and after every 50th iteration. The improved damping reads

$$a'_{\mathrm{new}} = \frac{a_{\mathrm{new}} + a_{\mathrm{old}}\mathrm{Damp}_{\mathrm{nu}}}{1 + \mathrm{Damp}_{\mathrm{nu}}}, \quad 0 \le \mathrm{damp}_{\mathrm{nu}} < \infty, \qquad (13)$$

where a_{new} calculated by Eq. (12) is substituted by a'_{new} taking into account the matrix from previous iteration, a_{old}, and the appropriate damping $\mathrm{damp}_{\mathrm{nu}}$. Considering the low-temperature range even improved damping fails to start the iteration process. In this case the solutions obtained at a higher temperature are used as initial values of a in Eq. (12) and then the temperature is gradually decreased by factor 0.85.

To obtain the magnetic-field distribution from a known vector potential distribution, the field $\mathbf{h}(\mathbf{r}) = \nabla \times \mathbf{A}(\mathbf{r})$ is calculated simultaneously for all $\mathbf{r}$ in the mesh using FFT. While transforming the real vector $\mathbf{A} = A_x\mathbf{e}_x + A_y\mathbf{e}_y$ to a complex number $A = A_x + iA_y$, we use the equation

$$\mathrm{FFT}^{-1}[(q_x - iq_y)\mathrm{FFT}(A_x + iA_y)] = (\nabla \times \mathbf{A}) \cdot \mathbf{e}_z - i\nabla \cdot \mathbf{A}, \qquad (14)$$

where FFT and FFT^{-1} denote direct and reverse two-dimensional fast Fourier transforms and q_x and q_y are the coordinates of the corresponding wave vector $\mathbf{q}_{ij} = q_x(i)\mathbf{e}_x + q_y(j)\mathbf{e}_y$ over all matrix cells.

After solving the Eilenberger equations the obtained magnetic-field distribution $h_E(\mathbf{r})$ is fitted with the London field distribution $h_L(\mathbf{r})$ finding the fitting parameter ξ_h. The normalized difference between these fields corresponding to $B=1$, $T=0.5$, and $\tau=1$ is shown in Fig. 1. The accuracy of the fitting exceeds 1%. Figure 2 shows the calculated $\xi_h(B,\tau)$ surface at $T=0.5$. The main effect on ξ_h arises from impurity scattering resulting in strong suppression of the vortex core size when τ decreases. Growing of ξ_h with T has been found previously[7] but with different shapes. These dependencies have different physical meanings: increasing of $\xi_h(\tau)$ is connected with approaching the clean limit, while growing of $\xi_h(T)$ is explained by divergency of the coherence length near T_c.

Figure 3 demonstrates the calculated field dependence of ξ_h in pure and dirty superconductors with different relaxation times τ and temperatures, (a) $T=0.2$ and (b) $T=0.5$. It is found that the shape of $\xi_h(B)$ depends strongly on temperature. At moderate temperatures (e.g., $T=0.5$) there is a minimum in $\xi_h(B)$. A minimum was found also in the order-parameter coherence length ξ_Δ (Ref. 6) (ξ_1 in notation of Ref. 1) reflecting interconnection between the magnetic coherence length $\xi_h(B)$ and the order-parameter distribution. Numerical calculations of $\xi_\Delta(B)$ show that embedding impurities results in suppression of the minimum in ξ_h and eventually leads to a monotonically decreasing function.[6] At low temperatures like $T=0.2$ the values of $\xi_h(B)$ monotonously increase when B is increased. At all temperatures, increasing of the scattering rate results in decreasing of ξ_h making the $\xi_h(B)$ curve more flat. As can be seen from Fig. 3 at $\tau \approx 0.5$ the magnetic-field dependence of ξ_h practically disappears. Suppression of ξ_h with increasing scattering rate is so strong that it excludes any possibilities of crossing the $\xi_h(B)$ curves at different τ. This behavior is different from the $\xi_\Delta(B)$ curves obtained in Ref. 6 where crossing of these curves is clearly visible. Absence of the crossing was also found in linearized approximation of the Eilenberger equation.[2] It probably reflects the importance of the length scale ξ_2 in calculation of ξ_h.

Minimum in $\xi_h(B)$ dependence was also found for clean d-wave superconductors.[17] Presence of nonmagnetic impurities could change the $\xi_h(B)$ dependence in d-wave superconductors because there are pair breakers decreasing the critical temperature.[18] As a result, nonzero density of state appears at

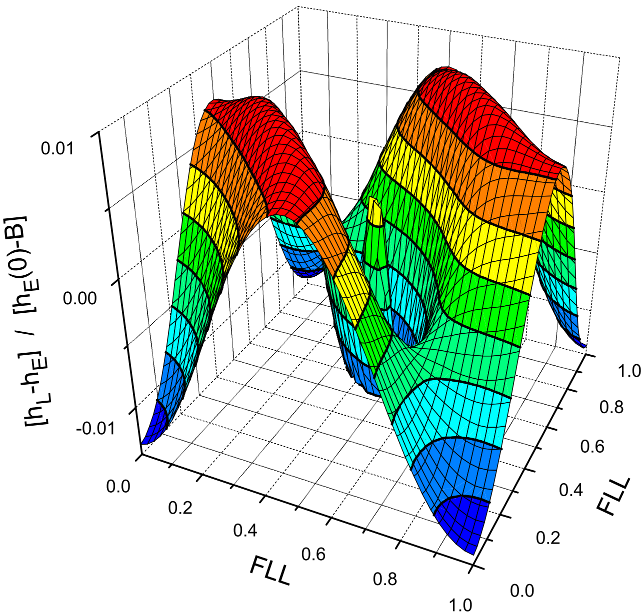

FIG. 1. (Color online) Normalized differences between the fields h_L and h_E calculated with the generalized London model and with the Eilenberger equation, respectively, for $B=1$, $T=0.5$, and $\tau=1$. The scales of the lengths are those of the flux-line lattice unit vectors.

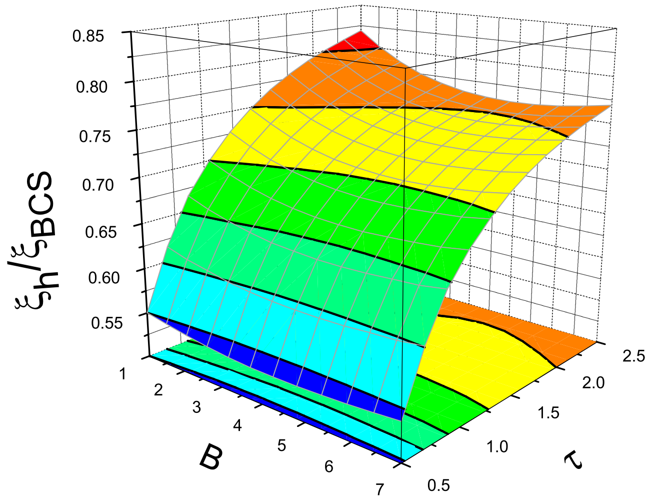

FIG. 2. (Color online) Three-dimensional plot of $\xi_h(B,\tau)$ at $T=0.5$.

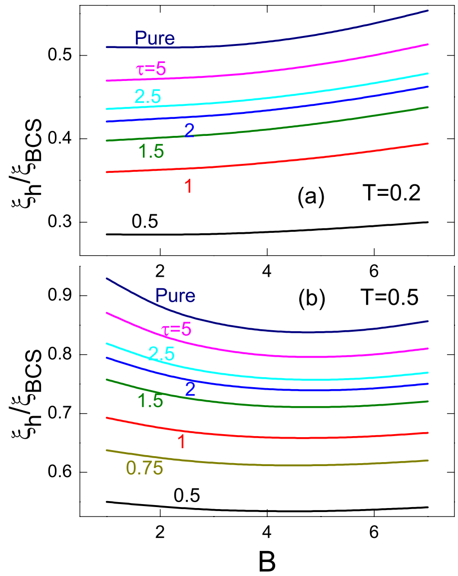

FIG. 3. (Color online) The calculated field dependence of ξ_h at different relaxation times τ and temperatures, (a) T=0.2 and (b) T=0.5.

small concentration of impurities even in the Meissner state. It could make the $\xi_h(B)$ dependence much more flat in comparison to clean d-wave superconductors similar to s-wave superconductors with strong scattering (see Fig. 3).

Figure 4 illustrates the behavior of $\xi_h(\tau)$ in a wide range of τ up to moderate scattering where τ is comparable with characteristic time τ_0 (to be defined below) at B=5 and T=0.2, 0.5, and 0.8. At small τ a sharp increase in $\xi_h(\tau)$ with increasing τ is evident. It resembles the prediction for $\xi_{\rm el}$ in dirty superconductors with uniform order parameter, where $\xi \propto \tau$.[11] In the clean limit we have $\tau \gg 1$ and asymptotically $\xi_h(\tau) - \xi_h(\infty) \propto 1/\tau$. The value of $\xi_h(5)$ is near the limit for clean superconductors $\xi_h(\infty)$, shown in Fig. 4 by the dotted lines.

Our numerical approach extends the theory of electromagnetic coherence length in the Meissner state[11] to the

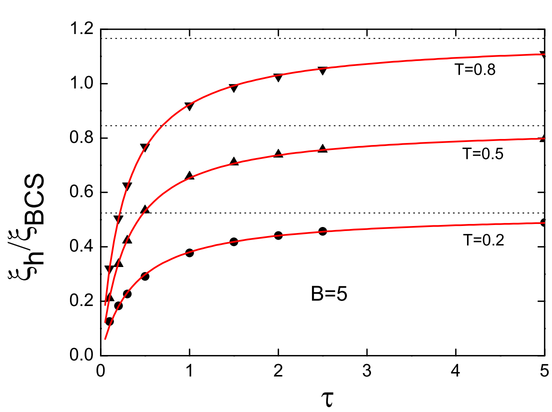

FIG. 4. (Color online) The dependence of ξ_h on impurity scattering at B=5 and T=0.2, 0.5, and 0.8. The solid lines represent fitting according to Eq. (18); the dotted lines denote the pure limit (τ=∞).

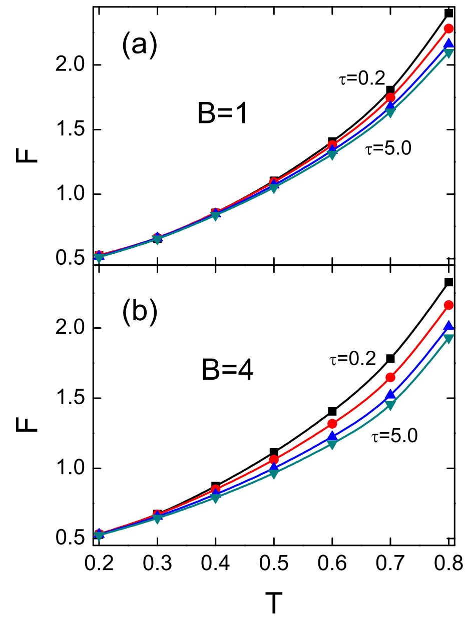

FIG. 5. (Color online) Temperature dependence of $F \equiv \xi_h/\xi_{\rm el}$ ratio at (a) B=1 and (b) B=4. From top to bottom the relaxation times τ are 0.2, 0.5, 1.5, and 5.0 in both plots.

mixed state. In our case the shape of the $\xi_h(\tau,T)$ dependence shows some differences which can be attributed to different definitions of the coherence length ξ_h in our model and $\xi_{\rm el}$.[11] The electromagnetic length $\xi_{\rm el}$ is defined by the relation[19]

$$\lim_{q\to\infty} qK(q,T)/K(0,T) = 3\pi/4\xi_{\rm el}, \tag{15}$$

where $K(q,T)$ can be calculated in terms of single-particle Green's functions. It governs the connection between the current and the vector potential in the Meissner state,

$$\mathbf{j}(\mathbf{q},T) = -\frac{c}{4\pi}K(\mathbf{q},T)\mathbf{A}(\mathbf{q}). \tag{16}$$

In framework of the Bardeen-Cooper-Schrieffer (BCS) theory we have

$$\xi_{\rm el}/\xi_{\rm BCS} = \frac{\pi\Delta}{2}\sum_{n\geq 0}\frac{1}{\tilde{\Delta}_n(1+u_n^2)^{3/2}}\Big/\sum_{n\geq 0}\frac{1}{1+u_n^2}. \tag{17}$$

Figure 5 shows the ratio $\xi_h/\xi_{\rm el}$ at (a) B=1 and (b) B=4 at different relaxation times τ=0.2, 0.5, 1.5, and 5.0. While uniform superconductors are considered in the calculation of $\xi_{\rm el}$, in our case ξ_h is calculated assuming a space-dependent self-consistent pairing potential in the vortex core. In the mixed state the Kramer-Pesch effect[5] is important at low temperatures resulting in difference of (Ref. 7) from $\xi_{\rm el}(T)$.[11] Intervortex interaction in the mixed state results in field dependence of ξ_h which is absent in $\xi_{\rm el}$. As can be seen from Fig. 5 the $\xi_h/\xi_{\rm el} \equiv F$ ratio is weakly dependent on relaxation time illustrating that both ξ_h and $\xi_{\rm el}$ have a similar dependence on τ. This demonstrates that characteristic relaxation times by impurities are similar in the Meissner and in the

TABLE I. The parameters ξ_{pure} and τ_0 used for fitting Eq. (18) at $B=5$ and $T=0.2$, 0.5, and 0.8.

T/T_c	0.2	0.5	0.8
ξ_{pure}/ξ_{BCS}	0.525	0.845	1.167
$\tau_0(2k_BT_c/\hbar)$	0.379	0.294	0.264

mixed state only with slight renormalization by the magnetic field in the last case. The main effect of nonuniformity of the order parameter is described by temperature dependence of ξ_{pure}. It means that the $\xi_h(B,T,\tau)$ dependence can be approximately attributed to two contributions: one is connected with core effects described by the $F(T)$ function (see Fig. 5) and the other is responsible for impurity scattering characterized by $\xi_{el}(\tau)$.

For comparison with experiments an analytical relation is needed. This approach was successfully used for description of the electromagnetic coherence length in the Meissner state.[11] To describe the mixed state we fit $\xi_h(\tau)$ in the considered range of impurity concentration by the function

$$\xi_h(B,T,\tau) = \frac{\xi_{pure}(B,T)}{1+\dfrac{\tau_0(B,T)}{\tau}}, \qquad (18)$$

where $\xi_{pure}(B,T)$ is the effective coherence length in clean superconductors.[7] The chosen $\xi_h(\tau)$ interpolation formula [Eq. (18)] has the same behavior as ξ_{el} in the limits of $\tau\to 0$ $\xi_h \propto \tau$ and at $\tau\to\infty$ $\xi_h\to\xi_{pure}$. The characteristic relaxation time $\tau_0(B,T)$ is given at temperatures $T=0.2$, 0.5, 0.8, and $B=5$ in Table I. As can be seen from Fig. 4, one-parameter fitting represents excellently the numerically calculated impurity dependence of $\xi_h(\tau)$ which has a similar shape as in the Meissner state[19] but with different temperature dependences of the fitting coefficients; in the Meissner state all parameters are decreasing functions of temperature but in our case $\tau_0(T)$ is a decreasing and $\xi_{pure}(T)$ an increasing function, resulting from core effects. Figure 6 demonstrates the $\tau_0(B,T)$ surface where decreasing of τ_0 with temperature holds out in a wide magnetic-field range. As can be seen from Fig. 5 the function F depends mainly on temperature. This means that characteristic relaxation time is similar in the Meissner and in the mixed states.

The model has been introduced to describe local properties of FLL.[7] Taking into account the nonlocal properties of FLL, an anisotropic nonlocal L_{ij} tensor can be used instead of London penetration depth.[20] This method is useful for description of the transition from triangular to square vortex lattice observed by SANS technique in high-T_c superconductors[21] and for explanation of the effects related to the fourfold anisotropy in field distribution in high-T_c superconductors.[20,22] The theory has been successfully applied also for interpretation of the results of μSR investigations.[23] Microscopic theory based on quasiclassical theory describing these effects has been developed,[17,24] and it was found that the nonlocal generalized London equation[16] model can be used as a reasonable approximation of the quasiclassical approach. A good test of the applicability of the model is calculation of the field distribution anisotropy of the single vortex in d-wave superconductors; the model predicts the sign changes in $H(r,\phi=0)-H(r,\phi=\pi/4)$ at some radius at low temperature, in agreement with the Eilenberger equations.[25] However a magnetic-field-dependent cutoff parameter is still needed for detailed description of the field distribution. The anisotropy of FLL is often unimportant for explanation of the magnetization data and transport measurements.[2] Thus, the more simple local model[7] can be used. It should be also noted that the shapes of the ξ_h/ξ_0 dependences on B are similar in the local and nonlocal models.[7,24] Using a field-dependent effective penetration depth, $\lambda(B)$, μSR measurements give information about anisotropy of the superconducting order parameter by measuring the slope of λ against magnetic field.[26] This method also requires an *a priori* defined cutoff parameter which cannot be determined directly from other experiments. It has been experimentally established for a variety of materials, which the value of λ_{ab} when $H\to 0$ agrees with the magnetic penetration depth measured by other techniques in the Meissner phase. Consequently, only $\lambda_{ab}(H\to 0)$ can be considered as true measure of the superconducting electron density.[27]

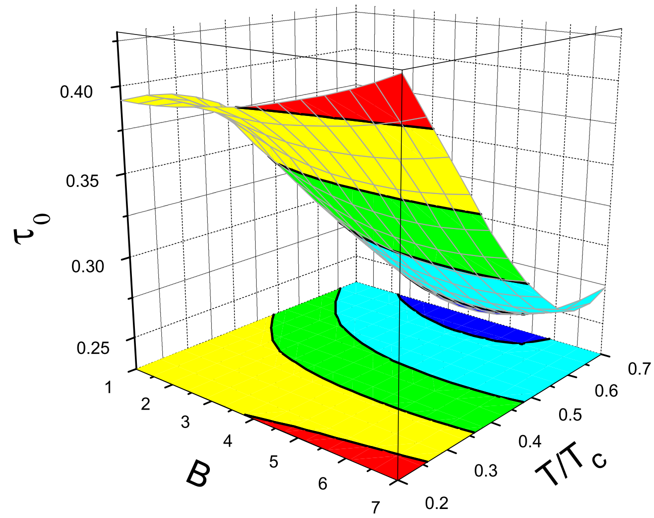

FIG. 6. (Color online) Three-dimensional plot of characteristic relaxation time $\tau_0(B,T)$.

To conclude, magnetic field, relaxation time, and temperature dependences of the effective coherence length in the mixed state are obtained by solving Eilenberger equations. This length determines the form factor of FLL. The minimum in the $\xi_h(B)$ dependence is found in superconductors at moderate temperatures. It will be interesting to check this prediction experimentally by SANS measurements. At low temperatures the value of $\xi_h(B)$ monotonously increases and the shape of the $\xi_h(B)$ curve is flattened along with decreasing scattering time τ (Fig. 3). In wide range of impurity concentrations the calculated results can be well expressed with a one-parameter fitting function.

This work was supported by the Wihuri Foundation, Finland.

*reilai@utu.fi
†Also at A. F. Ioffe Physico-Technical Institute, St. Petersburg 194021, Russia; miksaf@gmail.com

[1] J. E. Sonier, J. Phys.: Condens. Matter **16**, S4499 (2004).
[2] V. G. Kogan, R. Prozorov, S. L. Bud'ko, P. C. Canfield, J. R. Thompson, J. Karpinski, N. D. Zhigadlo, and P. Miranović, Phys. Rev. B **74**, 184521 (2006).
[3] K. O. Cheon, I. R. Fisher, V. G. Kogan, P. C. Canfield, P. Miranovic, and P. L. Gammel, Phys. Rev. B **58**, 6463 (1998).
[4] V. G. Kogan and N. V. Zhelezina, Phys. Rev. B **71**, 134505 (2005).
[5] L. Kramer and W. Pesch, Z. Phys. **269**, 59 (1974).
[6] P. Miranović, M. Ichioka, and K. Machida, Phys. Rev. B **70**, 104510 (2004).
[7] R. Laiho, M. Safonchik, and K. B. Traito, Phys. Rev. B **76**, 140501(R) (2007).
[8] Z. Hao, J. R. Clem, M. W. McElfresh, L. Civale, A. P. Malozemoff, and F. Holtzberg, Phys. Rev. B **43**, 2844 (1991).
[9] A. Yaouanc, P. Dalmas de Reotier, and E. H. Brandt, Phys. Rev. B **55**, 11107 (1997).
[10] T. Kita, Phys. Rev. B **68**, 184503 (2003).
[11] T. R. Lemberger, D. M. Ginsberg, and G. Rickayzen, Phys. Rev. B **18**, 6057 (1978).
[12] L. DeBeer-Schmitt, C. D. Dewhurst, B. W. Hoogenboom, C. Petrovic, and M. R. Eskildsen, Phys. Rev. Lett. **97**, 127001 (2006).
[13] A. A. Abricosov, L. P. Gor'kov, and I. E. Dzyaloshinskii, *Quantum Field Theoretical Methods in Statistical Physics* (Prentice-Hall, Englewood Cliffs, NJ/Pergamon, London, 1963).
[14] P. Miranović and K. Machida, Phys. Rev. B **67**, 092506 (2003).
[15] N. Schopohl and K. Maki, Phys. Rev. B **52**, 490 (1995).
[16] M. Safonchik, Ph.D. thesis, Turku University, 2007.
[17] R. Laiho, M. Safonchik, and K. B. Traito, Phys. Rev. B **73**, 024507 (2006).
[18] L. S. Borkowski and P. J. Hirschfeld, Phys. Rev. B **49**, 15404 (1994).
[19] M. Tinkham, *Introduction to Superconductivity* (McGraw-Hill, New York, 1975).
[20] M. Franz, I. Affleck, and M. H. S. Amin, Phys. Rev. Lett. **79**, 1555 (1997).
[21] S. P. Brown, D. Charalambous, E. C. Jones, E. M. Forgan, P. G. Kealey, A. Erb, and J. Kohlbrecher, Phys. Rev. Lett. **92**, 067004 (2004).
[22] I. Affleck, M. Franz, and M. H. Sharifzadeh Amin, Phys. Rev. B **55**, R704 (1997).
[23] M. H. S. Amin, M. Franz, and I. Affleck, Phys. Rev. Lett. **84**, 5864 (2000).
[24] R. Laiho, M. Safonchik, and K. B. Traito, Phys. Rev. B **75**, 174524 (2007).
[25] R. Laiho, E. Lähderanta, M. Safonchik, and K. B. Traito, Phys. Rev. B **71**, 024521 (2005).
[26] R. Kadono, S. Kuroiwa, J. Akimitsu, A. Koda, K. Ohishi, W. Higemoto, and S. Otani, Phys. Rev. B **76**, 094501 (2007).
[27] J. E. Sonier *et al.*, Phys. Rev. B **76**, 134518 (2007).