

# Thermal creep induced by cooling a superconducting vortex lattice

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A perturbed system relaxes towards an equilibrium given by a minimum in the potential-energy landscape. This often occurs by thermally activated jumps over metastable states. The corresponding dynamics is called creep and follows Arrhenius' law. Here we consider the situation where the equilibrium position itself depends on temperature. We show that this effect occurs in the vortex lattice of the anisotropic superconductor 2H-NbSe<sub>2</sub> when the magnetic field is tilted away from the principal axes, and that it leads to the peculiar appearance of creep when cooling the sample. Temperature determines the system's equilibrium state and at the same time brings the system back to equilibrium, playing a dual and antagonistic role. We expect that cooling-induced creep occurs in correlated systems with many degrees of freedom, allowing one to tune the equilibrium state via heat treatment.

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## I. INTRODUCTION

Superconducting vortices are lines of quantized magnetic flux  $\Phi_0 = hc/2e$  where the superconducting order parameter is depressed at a length scale of the order of the superconducting coherence length  $\xi$  and circular supercurrents are flowing at a length scale of the order of the penetration depth  $\lambda$ . Vortices arrange in a lattice, which can be triangular, square, or disordered due to the interactions of vortices with the crystalline environment [1–5]. Vortices tend to be pinned on material defects or inclusions consisting of places where superconductivity is depressed on length scales of the order of  $\xi$ . When varying the applied magnetic field, its strength, or direction, vortices enter or exit the sample [1–3]. This produces vortex motion, which is often counteracted by pinning. The action of the pinning landscape results in long-lived out-of-equilibrium vortex distributions that relax through thermal creep over a manifold of barriers [1–3, 6–10]. The phenomenon of creep has been observed in interacting systems of particles, such as colloids, polymers, solids consisting of mixtures, or in lattices of entities formed through electronic interactions (domain walls or skyrmions) [11–15]. To study pinning and creep vortex lattices are often considered as a model system [15–18]. Most efforts to understand vortex creep have focused on trying to immobilize vortices and thereby increase the critical current  $j_c$  for applications [19]. But the interaction between vortices

and the underlying superconducting material is very rich and can produce counterintuitive phenomena. Here we find that the equilibrium state towards which the system creeps can be modified with the temperature, leading to a phenomenon that we term *self-imposed creep*; see Fig. 1. We study vortex creep in the superconductor 2H-NbSe<sub>2</sub> and observe that after

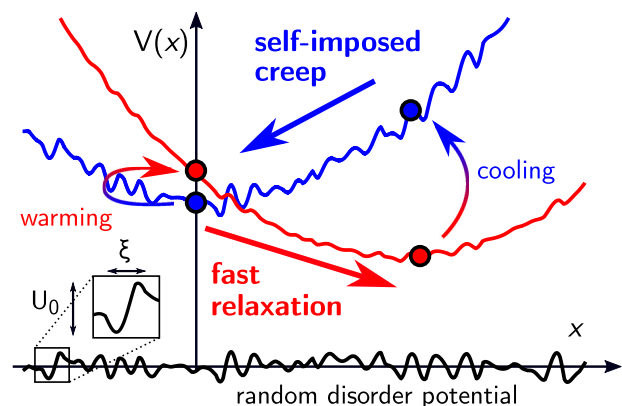


FIG. 1. Schematic response of a particle confined in a potential subject to random disorder. The disorder potential is characterized by local maxima of size  $U_0$  and over a length scale  $\xi$  (the superconducting coherence length in the case of vortices in superconductors). Upon warming, the particle rapidly moves to a newly defined minimum as the disorder is thermally smeared out [indicated by effectively smaller wells (red)]. When cooling, the potential changes and the particle now occupies a metastable state relative to the new global minimum. As the motion is impeded by pinning barriers (blue), the particle creeps by thermal relaxation towards the new minimum.

reaching the equilibrium state by warming, the creep motion of the vortex lattice reappears upon cooling. We capture this phenomenon within a model for creep in disordered systems.

The disappearance and reappearance of vortex creep in 2H-NbSe<sub>2</sub> is discussed and quantified in Sec. II. In Sec. III, we propose a generic model to describe the phenomenon of self-imposed creep. Section IV places in perspective the theoretical predictions with the experimental observations.

## II. VORTEX CREEP IN 2H-NbSe<sub>2</sub>

The layered compound 2H-NbSe<sub>2</sub> shows weak pinning and small creep rates [20,21]. Due to the material anisotropy, the properties of vortex creep depend on the field orientation relative to the crystallographic *c* axis. To quantify this effect, we tilt the field angle (after cooling the system in a field  $H = 0.85$  T) by  $\theta_H = 70^\circ$  away from the *c* axis.

2H-NbSe<sub>2</sub> is an anisotropic superconductor, with the upper critical field varying as a function of the field direction and quantified by the anisotropy  $\varepsilon \equiv H_{c2,ab}/H_{c2,c} \approx 1/3$ . This leads to a distortion of the vortex lattice, discussed in detail in Refs. [22–29], which can be neglected for our purposes. A rotation of the magnetic field to the field angle  $\theta_H$  induces creep, as the modification in the flux direction is counteracted by disorder present in the system. In isotropic systems, vortices would creep towards  $\theta_H$ ; see Fig. 2. However, in anisotropic systems, the superconducting anisotropy modifies the angle for the equilibrium configuration of the vortex lattice to  $\theta_B$ . Minimization of the free energy yields [2,30–32]

$$\sin(\theta_B - \theta_H) = \frac{H_{c1}}{H} \frac{(1 - \varepsilon^2) \sin \theta_B \cos \theta_B}{(\varepsilon^2 \sin^2 \theta_B + \cos^2 \theta_B)^{1/2}}, \quad (1)$$

up to a logarithmic correction of the order of unity. In 2H-NbSe<sub>2</sub>, using  $H_{c1}(T=0) \approx 200$  G for the lower critical field, we find  $\theta_B - \theta_H \approx 0.8^\circ$ , i.e., vortices are more inclined towards the *ab* plane than the external field; see Fig. 2(a). In the regime where pinning is weak, i.e., where the Bean length [33]  $\ell_B \approx cB/4\pi j_c$  is larger than the sample width  $w$  and thickness  $d \lesssim w$ , vortices in the critical state [33,34] are straight and oriented along the angle  $\theta_B^c = \theta_B - (w/2\ell_B) \sin \theta_B < \theta_B$ ; see Fig. 2(a). With  $j_c \sim 10^3$  A/cm<sup>2</sup>, this critical angle deviates from the equilibrium angle  $\theta_B$  as  $\theta_B - \theta_B^c \approx 0.5^\circ$ . We further notice that  $\theta_B(T)$  dominantly depends on the current distribution via the penetration depth  $\lambda(T)$  and is reflected in the above equation through  $H_{c1}(T)$ . For small temperature changes  $H_{c1}(T) = H_{c1}(T_0) + \delta H_{c1}$ , large fields  $H \gg H_{c1}$ , and a tilt angle  $\theta_B$  away from 0 or  $\pi/2$ , the angle changes to  $\theta_B(T) = \theta_B(T_0) - \delta\theta_B$ , with

$$\delta\theta_B \approx -[\theta_B(T_0) - \theta_H] \delta H_{c1}/H_{c1}(T_0). \quad (2)$$

The magnitude of this effect can be estimated using the Ginzburg-Landau scaling  $H_{c1}(T) \approx H_{c1}(0)(1 - T/T_c)$ , which provides a relative change in  $H_{c1}$  between the experiment's low ( $T_0 = 150$  mK) and high ( $T = 2$  K) temperatures of  $\delta H_{c1}/H_{c1} \approx -0.29$ , and hence  $\delta\theta_B \approx 0.3^\circ$ .

To investigate vortex creep in this system, we use a dilution-refrigerator scanning tunneling microscope (STM) in a three-axis vector magnet [35]. The cleaved 2H-NbSe<sub>2</sub> sample ( $\sim 1 \times 1 \times 0.2$  mm<sup>3</sup>)—grown with iodine vapor deposition—is studied with a gold tip, sharpened and cleaned

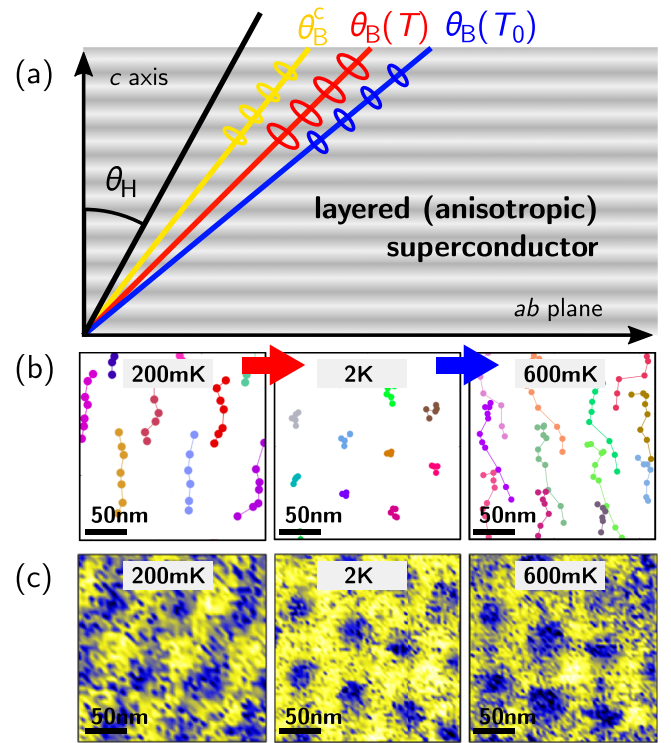


FIG. 2. (a) Schematic vortex alignment for an anisotropic superconductor in a tilted magnetic field ( $\theta_H$ , black). The  $T$ -dependent current patterns—indicated as rings—define different equilibrium orientations  $\theta_B$  at low ( $T_0$ , blue) and higher ( $T > T_0$ , red) temperature. When tilting the magnetic field to  $\theta_H$ , vortices creep towards the equilibrium angle  $\theta_B^c < \theta_B(T_0)$  (yellow). (b) Vortex positions extracted from individual STM images (points) recorded at fixed temperature. Each vortex (core size  $\xi \approx 7$  nm) is identified by a color and the lines join the position between two consecutive images. At low  $T$ , out-of-equilibrium vortices move between subsequent frames. The motion stops upon warming to 2 K. At this temperature, vortices are fixed in a lattice. When cooling, motion reappears again and the average image shows that vortices do not stay at the same position in subsequent steps. (c) Time average of the series of STM images used in (b).

*in situ* [36]. At the field strength  $H = 0.85$  T ( $H \parallel c$ ), the magnetization is practically reversible so that no vortex motion is measured within days. Vortex motion only appears when rotating the field towards the layers (after field cooling to the base temperature  $T_0 = 150$  mK with the magnetic field directed along the *c* axis).

While specific vortex positions can be identified from an individual STM image (taking 23 minutes to record), consecutive images allow one to measure vortex motion over time. Figures 2(b) and 2(c) show the evolution of vortex positions at three fixed temperatures and the corresponding time-averaged STM image, respectively. It highlights the presence of vortex creep at low temperature, its disappearance upon warming, and its reappearance upon cooling. Note that vortices always move along the direction of the component of the magnetic field within the layers. In agreement with earlier reports [20,37–39], we observe a weak modulation of the vortex velocity when vortices travel by a distance of the order of the intervortex separation. The overall vortex displacement

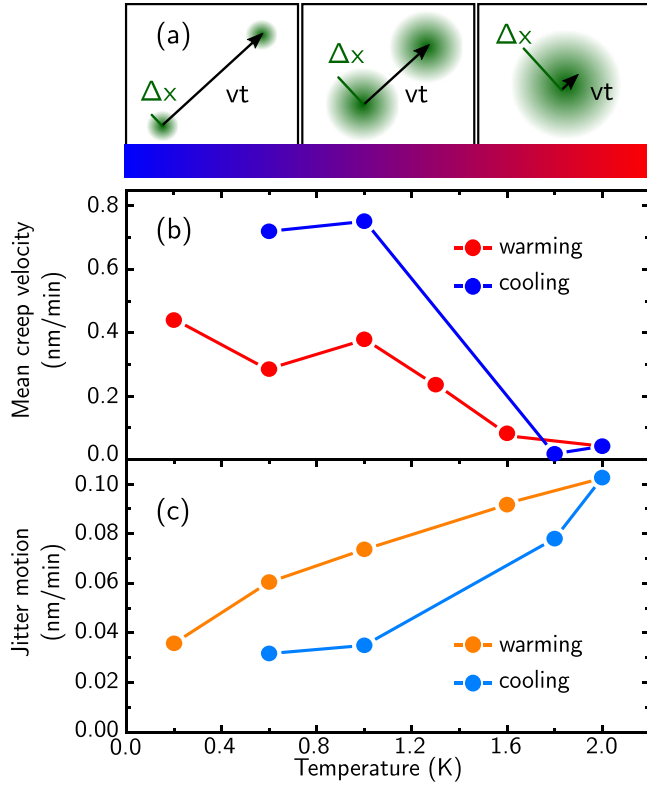


FIG. 3. (a) Schematic mean creep velocity (black arrow) and jitter motion (green cloud) for different temperatures. (b) Average creep velocity of vortices in 2H-NbSe<sub>2</sub> at tilted field and fixed temperature, recorded on a warming (red) and cooling (blue). (c) Standard deviation of the vortex jitter motion upon warming (orange) and cooling (light blue). The behavior agrees with the phenomenology of self-imposed creep.

translates to an average angular velocity  $\dot{\theta}_b \approx 10^{-3}$  deg/h of vortices tilting to new equilibrium positions [see Fig. 2(a)] and the apparent absence of a measurable decay time within the observation window ( $\sim 5$  h) is consistent with the lower bound  $|\theta_b - \theta_b^c|/\dot{\theta}_b \gtrsim 100$  h provided by the angular misalignment discussed above.

Without reinitializing the vortex state, imaging is repeated at different temperatures, while the system is kept at finite field strength and angle. Thus, temperature determines the preparation of the vortex state, as in Refs. [40–42]. Figure 3(a) schematically shows how the vortex displacement and jitter motion evolve upon changing temperature. In Figs. 3(b) and 3(c), the evolution of the average creep velocity  $v(T)$  and jitter motion  $\Delta x(T)$  is shown for a warming (red) and cooling (blue) cycle. Each data point is obtained from a series of STM frames. To find  $v(T)$ , we determine the position  $\mathbf{r}_j^i$  of vortex  $j$  in frame  $i$  and evaluate the average displacement for each vortex per frame, given by  $\delta \mathbf{r}_j = |\mathbf{r}_j^{n_j} - \mathbf{r}_j^1|/(n_j - 1)$ , where  $n_j$  denotes the number of frames where the  $j$ th vortex appears. Averaging over all  $N_v$  vortices for a given temperature, we arrive at the average creep velocity  $v(T) = (t_f N_v)^{-1} \sum_{j=1}^{N_v} \delta \mathbf{r}_j$ , with  $t_f$  the time for measuring one frame. To quantify the jitter motion, we evaluate the average jitter displacement  $\delta s_j = (\frac{1}{n_j - 1} \sum_{i=2}^{n_j} |\delta \mathbf{r}_j^i|) - \delta \mathbf{r}_j$ , with  $|\delta \mathbf{r}_j^i|$  the vortex displacement

between two subsequent frames  $i-1$  and  $i$ . The average over all frames at a fixed temperature now provides the jitter motion,  $\Delta x(T) = N_v^{-1} \sum_{j=1}^{N_v} \delta s_j$ .

The average creep velocity decreases upon warming and vanishes above 2 K. Upon cooling, however, a finite velocity reappears. If the vortices were to reach a temperature-independent minimum upon warming, the jitter motion would decrease upon cooling without a reappearance of creep motion. The reversible directed vortex motion upon thermal cycling is therefore a clear signature of self-imposed creep.

### III. MODEL FOR SELF-IMPOSED CREEP

To capture the main observations, we discuss a simple, yet quite generic model for creep in disordered landscapes. First, let us note that motion is driven by thermal fluctuations, through an Arrhenius-type activation process across the pinning barriers. Then, the timescale  $t \sim \tau \exp(U_0/k_B T)$  for thermal activation is determined by (i) the temperature  $T$ , (ii) the energy barrier  $U_0$ , and (iii) a microscopic timescale  $\tau = \omega^{-1}$  ( $\omega$  is an attempt frequency to overcome the barrier). Creep is observed when the timescale of the experiment is such that  $U_0 \sim k_B T \ln(\omega t)$ . To see the consequences of modifying the equilibrium, let us consider the problem of a particle confined in a one-dimensional parabolic trap  $V_0(x) = kx^2/2$  and subject to a driving force  $V_d(x) = -Fx$ . To map the results of our simple model to a many-body system, we associate  $x$  to a macroscopic observable, such as the position (or angle) of the vortex lattice for a given field direction. The force produces a drive towards the equilibrium position  $\bar{x} = F/k$ . In addition, we consider a disorder landscape  $V_p(x)$  characterized by a typical depth  $U_0$  and width  $\xi$ , with  $k\xi^2/U_0 \ll 1$ ; see Fig. 1. We approximate the bare potential between neighboring minima  $x_{\pm} = x \pm \xi$  by  $V_p(x + \delta) \approx [1 - (\delta/\xi)^2]U_0$ . The overall potential  $V(x) = V_0 + V_d + V_p$  features local minima in the range  $\bar{x} - 2U_0/k\xi < x < \bar{x} + 2U_0/k\xi$ , and the position-dependent activation barrier  $U_b(x)$  to move from  $x_-$  to  $x_+$  (we assume  $x < \bar{x}$ ) is

$$U_b(x) = [(x - \bar{x})k\xi]^2/4U_0 + U_0 + (x - \bar{x})k\xi. \quad (3)$$

The thermally activated motion in the opposite direction, i.e., from  $x_+$  to  $x_-$ , is penalized by an additional energy  $-2(x - \bar{x})k\xi > 0$ . A particle initially far from the minimum  $\bar{x}$  will glide down the potential until reaching  $\bar{x} - 2U_0/k\xi$  from where it will be thermally activated across ever-growing barriers. After a time  $t$ , the particle has reached a position  $x_T$  satisfying Arrhenius' condition

$$U_b(x_T) = k_B T \ln(\omega t), \quad (4)$$

as smaller barriers have been overcome in exponentially shorter activation times. Inserting Eq. (3) into the condition (4), we obtain

$$x_T = \bar{x} - (2U_0/k\xi)[1 - \sqrt{(k_B T/U_0) \ln(\omega t)}]. \quad (5)$$

If  $x_T$  is still far from  $\bar{x}$  in the sense  $\bar{x} - x_T \gg k_B T \ln(\omega t)/k\xi$  [translating to  $T \ll T_b \equiv (U_0/k_B)/\ln(\omega t)$ ], the particle moves with an average velocity  $v \approx 2\xi/t$ . If, however, the particle has relaxed in the vicinity of the global minimum, the thermal activation becomes almost equally probable in both directions.



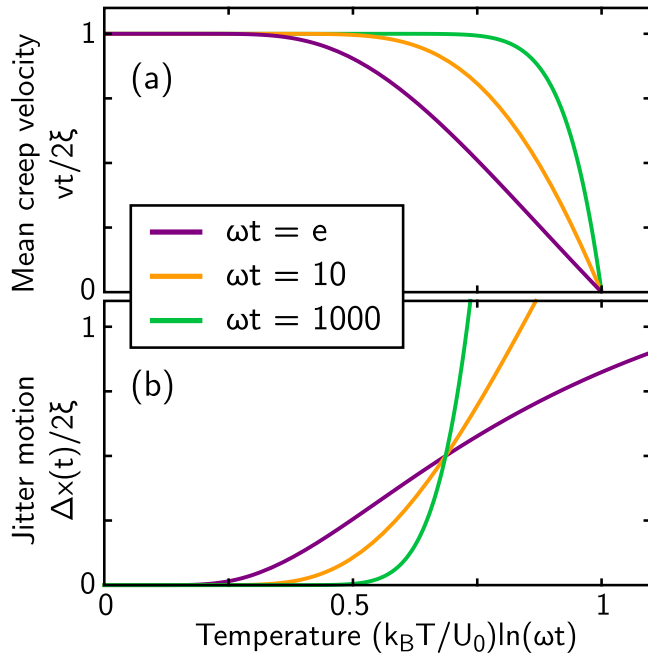


FIG. 4. (a) Average creep velocity  $v$  (in units of  $2\xi/t$ ) as a function of temperature [see Eq. (6)] for different timescales  $\omega t$ . When temperature affects both the activation dynamics and the global potential minimum (see Fig. 1), the velocity profile is traced reversibly upon warming and cooling. (b) Standard deviation  $\Delta x$  of the mean particle displacement, quantifying the isotropic thermal motion, or jitter; see Eq. (7).

Accounting for this bidirectional, yet asymmetric motion, we find a net average creep velocity

$$v \approx \frac{2\xi}{t} \left( 1 - \exp \left\{ -\frac{4U_0}{k_B T} [1 - \sqrt{k_B T \ln(\omega t)/U_0}] \right\} \right). \quad (6)$$

The temperature dependence of  $v$  is shown in Fig. 4(a) for different values  $\omega t$ . It is interesting to note that although the thermal activation energy increases upon warming, the creep velocity decreases. The validity of the above result is limited to temperatures  $T < T_b$ . For larger temperatures, the disorder landscape becomes irrelevant, as the particle relaxes to the global minimum within a time  $\sim \omega^{-1}$ .

The directed relaxation via creep [Eq. (6)] is blurred by an isotropic contribution, or jitter motion with zero mean and amplitude (standard deviation) given by

$$\Delta x \sim \xi \sqrt{\omega t} \exp \left\{ -\frac{2U_0}{k_B T} [1 - \sqrt{k_B T \ln(\omega t)/U_0}] \right\}. \quad (7)$$

This result is obtained by describing the (forward-backward) activation across the barriers as a stochastic process (random-walk motion), where the variance  $(\Delta x)^2 = \xi^2 \omega_{rw} t$  of the displacement grows linearly in time and is determined by the random-walk attempt frequency  $\omega_{rw} = \omega \exp[2(x_T - \bar{x})k\xi/k_B T]$ . The exponent only depends on the difference between the barrier heights for forward and backward motion. This jitter motion persists beyond the disappearance of creep motion and reaches a similar magnitude when  $k_B T \ln(\omega t) \sim U_0$ , as shown in Fig. 4(b).

When changing temperature, one has two distinguishable cases: one where the global minimum  $\bar{x}$  is constant, and one where  $\bar{x}(T)$  depends on  $T$  through a temperature-dependent force  $F(T)$ . In the first case, the average creep velocity is given by the local disorder landscape seen by the particle and hence follows Eq. (6). For  $T > T_b$ , the system is fully relaxed and the particle reaches  $\bar{x}$ . Meanwhile, the magnitude of the jitter motion continuously increases; see Fig. 4(b). Subsequent cooling lowers the thermal energy and the particle's motion freezes in place at  $\bar{x}$ . This is the conventional response expected for creep. In the second case, when  $\bar{x}$  depends on  $T$  through a temperature-dependent force  $F(T)$ , the situation is drastically different. A temperature change  $\Delta T$  imposes a shift  $\bar{x}(T + \Delta T) - \bar{x}(T) \gg \xi$  of the global equilibrium position. Upon warming, the velocity and jitter motion are similar to the previous case. However, the values of  $x(t)$  are different and, upon cooling, triggers particle motion tracing back (in magnitude) the warming curve, as schematically shown in Fig. 1. The appearance of creep upon cooling is caused by the two antagonistic roles of temperature—defining a  $T$ -dependent minimum and dictating the thermal relaxation towards it—and presents the most prominent manifestation of self-induced creep.

#### IV. DISCUSSION

We compare our observations in 2H-NbSe<sub>2</sub> with the above model. The typical barrier to overcome during pinning by thermal fluctuations is given by Arrhenius' law  $U_0 = k_B T \ln(\omega t)$ . To observe both the equilibrium phase at high temperature and reentrant creep at low temperature, it is important that the temperature of the experiment is of the order of  $(U_0/k_B)/\ln(\omega t)$ . In contrast to the pinning energy of one defect site, here  $U_0$  denotes the energy barrier for vortex creep [43,44]. Weak collective pinning theory [1,2,45] provides the estimate  $U_0 \sim k_B T_c [(j_c/j_{dp})(B/H_{c1})^3/Gi]^{1/2} v$ , with the depairing current  $j_{dp} = c\Phi_0/12\sqrt{3}\pi^2\lambda^2\xi$ , the Ginzburg-Levanyuk [46,47] number  $Gi \sim [T_c/H_c(0)^2\xi^3]^2$ , the condensation energy  $H_c(0)^2\xi^3 = \Phi_0^2\xi/8\pi^2\lambda^2$ , and  $v \sim 10^{-3}$  a small numerical number. From Refs. [48–53], we infer  $j_c/j_{dp} \sim 10^{-6}$  and  $Gi \sim 8 \times 10^{-7}$  and obtain  $U_0 \sim 5k_B T_c$ . This order-of-magnitude estimate is compatible with  $U_0 \lesssim k_B T \ln(\omega t)$ , provided  $\omega t \approx 10^3$ – $10^4$ , which is slightly larger than the values considered above (Fig. 4). Finally, the temperature-dependent current pattern in the layered superconductor naturally provides the necessary force  $F(T)$  to induce creep upon cooling a tilted vortex lattice. Given the simplicity of the one-dimensional model, the agreement is remarkable. All important features predicted by the model—the disappearance and reappearance of the directed motion, together with the temperature evolution of the jitter motion [behaviors of  $v(T)$  and  $\Delta x(T)$ , shown in Figs. 3(b) and 3(c) and in Figs. 4(a) and 4(b)]—are found in the experiment.

Given that the experimental timescale spans several minutes, our observation  $\omega t \approx 10^3$ – $10^4$  suggest a value for  $\omega$  of the order of one Hz. While a route for accurate determination of this attempt frequency is still lacking, the estimate  $\omega = \alpha_L/\eta$  for a single vortex depends on the vortex viscosity  $\eta$  and on the Labusch parameter [54]  $\alpha_L$  (which in turn

relates to the averaged potential curvature [55]). An analysis assuming vibrations with large  $k$  vectors reports frequency values in the range  $10^6$ – $10^{10}$  Hz [56]. In our case however, vortices are not isolated, but rather interact nonlocally with many vortices [2,3,57–59]. Low  $k$  vectors, or wavelengths comparable to the sample size, lead to highly dispersive elastic moduli which modify the attempt frequency by orders of magnitude [58–60]. Similar to our observation, previous measurements of slow vortex dynamics have reported [61,62] very low-frequency values for thermal motion and creep. Creep rates observed in layered cuprate superconductors involve extremely large timescales, indicating the relevance of collective creep [2,3,58,60]. Thus, even if the attempt rate of individual vortices is large, the dynamics as a lattice involves rates that are many orders of magnitude smaller. The temperature is far from melting, thus favoring collective rather than a single-vortex dynamics [63,64]. It is this near-equilibrium configuration with ultrasmall collective dynamics that allows for the observed cooling-imposed creep in our experiments.

The creep discussed here is very slow and the vortex velocity does not decay within our experimental time. However, the creep rate  $S = d \ln(j)/d \ln(t)$  can assume a sizeable value compatible with the suggested lower bound [21]  $S > (T/T_c)Gi^{1/2}$ . Actually, 2H-NbSe<sub>2</sub> is among the materials with the lowest creep rates, close to MgB<sub>2</sub> [21]. Creep between metastable vortex states that occur near the order-disorder transition of the vortex lattice in 2H-NbSe<sub>2</sub> or related to domain formation of lattices with different orientations in MgB<sub>2</sub> has been reported [65–67]. Motion then appears when modifying the relative strength of competing interactions, and it might well occur that the equilibrium configuration at some particular locations is influenced by temperature. Collective motion is also found in the stochastic behavior of particle arrangements [11]. Depending on particle interactions, the dynamics transits from individual random motion to flocking. The timescale related to flocking motion shows a divergent behavior with increasing interaction. Other long-term dynamical behavior should appear in thermal effects and might

lead to self-imposed creep whenever there are two or more parameters influencing the behavior of the system. Mixtures, such as alloys, concrete or rocks [68], liquids resolidifying under stress [12], steel under stress [13], colloidal stems, and magnetic domain walls or skyrmions [15] are cases of complex systems where self-imposed creep may be induced from temperature-dependent interactions.

The model of self-imposed creep explains the critical state dynamics in 2H-NbSe<sub>2</sub> at tilted magnetic fields; in particular, the commonly unexpected appearance of vortex motion when cooling. Likely, the balanced thermal activation dynamics and the temperature-dependent equilibrium could be matched in other complex systems.

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