A system perturbed out of equilibrium often relaxes following Arrhenius law, e.g., through thermally activated motion (creep) over local maxima $U_b$ in the potential energy landscape. When $U_b$ depends on temperature and $T \sim T_b \sim U_b/k_B$, a change in temperature varies both the system’s equilibrium state and the relaxation dynamics. We study vortex lattices out of equilibrium in superconducting 2H-NbSe$_2$ relaxing over several days and show that thermally activated motion can appear when cooling the vortex lattice. We show that temperature determines system’s ground state and at the same time brings the system back to equilibrium, playing a dual and antagonistic role.

Magnetic field penetrates type-II superconductors in the form of quantized vortices which order in a triangular lattice$^1$. When varying the applied magnetic field, its strength or direction, vortices enter or exit the sample$^{2,3}$. Their motion is hindered by pinning at material defects, that can produce long-lived out of equilibrium distributions of vortices. The vortex lattice relaxes towards equilibrium by thermal activation over a potential landscape with a manifold of barriers often spanning many orders of magnitude in energy$^{2-9}$. In a uniaxial superconductor when the magnetic field is tilted away from the principal axes, circulating currents around vortices preferably flow within the layers due to the crystalline anisotropy. This establishes a complex current distribution and leads to a slight misalignment between the vortex lines and the applied magnetic field$^{10-14}$. The misalignment is temperature dependent, because the current distribution is related to the penetration depth $\lambda(T)$, providing a situation which is quite unique—temperature simultaneously plays the role of a parameter that modifies equilibrium and that drives the system back to equilibrium after a perturbation through creep.

Quite generally, a system’s equilibrium state depends on its temperature $T$. Upon changing this thermodynamic parameter, the system migrates towards a new equilibrium following an energy gradient (or force). If the system can do so without obstacles, relaxation occurs within microscopic timescale $\tau = \omega^{-1}$, dictated by the attempt frequency $\omega$ to move to a new equilibrium. On the contrary, if reaching the new minimal energy state requires to overcome an energy barrier $U_b$ relaxation occurs through thermal fluctuations and there is an Arrhenius-type activation process on the timescale $t_\tau \sim \tau \exp(U_b/k_B T) \gg \tau$. The temperature is related to barrier height, attempt frequency and experimental timescale $t$ through $T \sim (U_b/k_B)/\ln(\omega t)$, which is also the temperature below which creep is observed. For high barriers $U_b \gg k_B T$, the system remains locked in a metastable state on the timescale $t \ll t_\tau$. For moderately high barriers, $t_\tau \lesssim t$ and thermal activation leads to a measurable relaxation. It is this dual and antagonistic role of the temperature (i) defining the thermodynamic equilibrium and (ii) dictating the rate of thermal motion, which is essential to the new creep effect discussed below. We therefore term this phenomenon self-imposed creep. Following our theoretical discussion, we provide experimental evidence for the observation of self-imposed creep for vortices in the anisotropic superconductor $2H$-NbSe$_2$.

Let us start by considering the problem of a particle in one dimension confined in a parabolic trap $V_0(x) = kx^2/2$ and subject to a tilting force $V_t(x) = -Fx$. The
force produces a drive towards the equilibrium position \( \dot{x} = F/k \). In addition, we introduce a disorder landscape \( V_p(x) \) characterized by typical depth \( U_0 \) and width \( \xi \), with \( k\xi^2/U_0 \ll 1 \), see Fig. 1. We approximate the barriers separating neighboring minima by \( U_0[1 - (x/\xi)^2] \). The overall potential \( V(x) = V_0(x) + V_d(x) + V_p(x) \) features local minima in the range \( x - 2U_0/k\xi < x < x + 2U_0/k\xi \).

For this potential, the activation energy \( U_b(x) \) to moving from one local minimum at \( x - \xi \) to the next lower one at \( x + \xi \) (with \( x < \bar{x} \)) is

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

(1)

The thermally activated motion in the opposite direction, i.e. from \( x + \xi \) to \( x - \xi \) 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. After a time \( t \), the particle has reached a position \( x_T \) satisfying Arrhenius’ condition

\[
U_b(x_T) = k_n T \ln(\omega t).
\]  

(2)

Inserting Eq. (1) into Eq. (2) and solving for \( x_T \) we find

\[
x_T = \bar{x} - (2U_0/k\xi)\left[1 - \sqrt{(k_n T/U_0) \ln(\omega t)}\right]
\]  

(3)

If \( x_T \) is still far from \( \bar{x} \), in the sense \( (\bar{x} - x_T)/k\xi \gg k_n T \), the particle moves with an average velocity \( v \sim 2\xi \omega \exp(-U_b/k_n T) \sim 2\xi/\omega t \). On the other hand if the particle has relaxed in the vicinity of the global minimum, the thermal activation becomes almost equally probable in both directions. This bidirectional, yet asymmetric, motion provides an average creep velocity

\[
v \sim \frac{2\xi}{t} \left[1 - \exp\left(-\frac{4U_0}{k_n T} \left(1 - \sqrt{\frac{k_n T \ln(\omega t)}{U_0}}\right)\right)\right].
\]  

(4)

temperature-dependence of which is shown in Fig. 2(a), for different values \( \omega t \). The validity of this result is limited to temperatures \( T < T_b \equiv (U_0/k_n) / \ln(\omega t) \). For a larger temperature, the disorder landscape becomes (essentially) irrelevant, as the particle relaxes within a time scale dictated by \( \omega^{-1} \). Close to the potential minimum, the thermal activation in the forward and backward direction become almost equally probable. As a result the creep motion (4) is blurred by an isotropic contribution (jitter motion) with zero mean and standard deviation

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

(5)

This result is obtained by assuming a stochastic process yielding a random-walk motion, where the variance \( \langle \Delta x^2 \rangle = \xi^2 \omega_{rw} t \) of the displacement grows linearly in time and is determined by the random-walk attempt frequency \( \omega_{rw} = \omega \exp\left[2(\bar{x} - \bar{x})k\xi/k_n T\right] \). This jitter motion persists beyond the disappearance of creep motion and reaches a similar magnitude when \( k_n T \ln(\omega t) \sim U_0 \), as shown in Fig. 2(b).

When considering the temperature dependence, we have to distinguish two scenarios: 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. (4). 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. 2(c). Subsequent cooling lowers the thermal energy and the particle’s motion freezes in place at \( \bar{x} \). Instead, when \( \bar{x} \) depends on \( T \) through a temperature-dependent force \( F(T) \), the situation is radically different. \( \Delta T \) imposes a shift \( \bar{x}(T + \Delta T) - \bar{x}(T) \approx \xi \) of the global equilibrium position. The velocity profile \( v(T) \) looks similar to the pre-
have been discussed in detail in Refs. [17–24] and are related to the uniaxial anisotropy, \( \varepsilon \equiv H_{c2, ab}/H_{c2, c} \approx 1/3 \). The direction of the induction \( \mathbf{B} \) differs from the direction of \( \mathbf{H} \) by an angle \( \theta_n - \theta_a \), with \( \theta_n \) and \( \theta_a \) being the angle of the field orientation for \( \mathbf{H} \) and \( \mathbf{B} \) relative to the crystallographic \( c \) axis. Minimizing the free energy with respect to \( \theta_n \) for fixed \( \mathbf{H} \) and \( \theta_a \) one finds

\[
\sin(\theta_n - \theta_a) = \frac{H_{c1}}{H} \left( 1 - \varepsilon^2 \right) \sin \theta_n \cos \theta_n \\
\left( \varepsilon^2 \sin^2 \theta_n + \cos^2 \theta_n \right)^{1/2}
\]

(6)

up to a logarithmic correction of order unity.\(^{2,25-27}\) The temperature dependence of the equilibrium angle \( \theta_n \), which depends on the current distribution and on \( \lambda(T) \), is encoded in \( H_{c1}(T) \). For small changes of \( H_{c1}(T) = H_{c1}(T_0) + \delta H_{c1} \), large fields \( H \gg H_{c1} \) and a tilt angle \( \theta_a \) away from 0 or \( \pi/2 \) around \( T = T_0 \), the angle changes to \( \theta_n(T) = \theta_n(T_0) - \delta \theta_n \), with

\[
\delta \theta_n \approx -[\theta_n(T_0) - \theta_a] \delta H_{c1}/H_{c1}(T_0).
\]

(7)

In the regime where pinning is weak, i.e., where the Bean length\(^{28,29}\) \( \ell = cB/4\pi j_c \) (\( j_c \) is the critical current and \( c \) the speed of light) is larger than the sample width \( w \) and thickness \( d \lesssim w \), vortices in this critical state\(^{24,26}\) are straight and oriented along the critical angle \( \theta_n^{\circ} = \theta_n - (w/2d) \) \( \sin \theta_n < \theta_n \), see Fig. 3(a). With \( j_c \approx 10^3 \, \text{A/cm}^2 \) \( (w \approx 1 \, \text{mm}) \) the critical angle \( \theta_n^{\circ} \) deviates from the equilibrium angle \( \theta_n \) as \( \theta_n - \theta_n^{\circ} \approx \approx 0.5^\circ \). Using \( H_{c1}(T = 0) \approx 200 \, \text{G} \) for the lower critical field and \( \varepsilon \approx 1/3 \), we find \( \theta_n - \theta_n^{\circ} \approx 0.8^\circ \), i.e., vortices are more inclined towards the \( ab \)-plane than the external field, Fig. 3(a). Finally, the Ginzburg-Landau scaling \( H_{c1}(T) \approx H_{c1}(0)(1 - T/T_c) \), provides a relative change in \( H_{c1} \) between the experiment’s low \( (T_0 = 150 \, \text{mK}) \) and high \( (T = 2 \, \text{K}) \) temperatures of \( \delta H_{c1}/H_{c1} \approx -2/7 \). This gives \( \delta \theta_n \approx 0.3^\circ \).

To measure a tiny vortex motion we make consecutive STM images as a function of time (each one taken in 23 minutes, an image is shown in Fig. 3(b)). We extract the vortex displacements from the image series [Fig. 3(b)]. From experiments at different azimuthal angles, we infer that vortices always move along the tilt of the magnetic field. The vortex motion shows a weak modulation at distances which correspond to multiples of the inter-vortex distance. This self-matching effect has been reported earlier in 2H-NbSe\(_2\) and in disordered thin films\(^{30-33}\) and evidences that the lattice moves as a whole in the creep regime. Other than that, no decay of vortex creep velocity is observed within the overall measurement time \( \approx 5 \, \text{h} \). The vortex displacement \( (160 \, \text{nm}) \) during this time translates to an angular velocity \( \approx 10^{-3} \, \text{deg/h} \) of vortices tilting to new equilibrium positions, see Fig. 3(a). The large time scale for thermal decay is consistent with the misalignment between magnetic field and vortex axis discussed above, \( |\theta_n - \theta_n^{\circ}|/\theta_a \approx 100 \, \text{h} \). While creep is very slow and no decay in the vortex velocity is observed within our experimental time, the creep rate \( S = d\ln(j_c)/d\ln(t) \) may

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**FIG. 3.** (a) Schematics of vortex alignment for anisotropic superconductors in a tilted magnetic field (\( \theta_n \), black). Temperature-dependent current patterns—indicated as rings—define different equilibrium orientations \( \theta_n \) at low \( (T_0 \), blue) and higher \( (T > T_0 \), red) temperature. If the tilted field is installed by rotating it away from the \( z \)-axis to \( \theta_n \), vortices get pinned in a critical state \( \theta_n^{\circ} < \theta_n(T_0) \), yellow) with a non-equilibrium current pattern. (b) Average over a series of STM images taken at three fixed temperatures. While at low temperature vortices move between subsequent frames (blurring), the motion stops upon warming to 2 K and reappears when cooling again. (c) The vortex positions extracted from the series of STM images (23 minutes per frame).
still assume a seizable value compatible with the suggested lower bound $S > (T/T_c)G_i^{1/2}$. $G_i$ denotes the Ginzburg-Levanyuk number. In our case, the evaluation of the creep rate and its comparison with magnetization measurements is further complicated by the different form of critical state, i.e., torque versus field-gradient.

Imaging is repeated at different temperatures. It is important to stress here, that no reinitialization occurs. Rather the system is kept at finite field strength and angle, and solely the temperature is swept. In Fig. 4(a) and 4(b) we show the temperature dependence of the average creep velocity $v(T)$ and the jitter motion $\Delta x(T)$ for a set of vortices, respectively. Each data point is obtained from a series of STM frames. To find $v(T)$ we determine the position $r_{ij}^T$ of vortex $j$ in frame $i$ and evaluate the mean displacement for each vortex per frame, given by $\delta r_j = |r_{ij}^T - r_{ij}^0|/(n_{ij} - 1)$, where $n_{ij}$ 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 $v(T) = \frac{1}{N_v} \sum_{j=1}^{N_v} \delta r_j$ with $t_l = 23\text{min}$ the time for measuring one frame. To quantify the jitter motion, we evaluate the average jitter displacement $\delta s_j = \langle |r_{ij}^0 - r_{ij}^1| \rangle - \delta r_j$, with $\langle |r_{ij}^0| \rangle$ the vortex displacement between two subsequent frames $i-1$ and $i$. The average over all frames at a fixed temperature now provides the jitter displacement $\Delta s(T) = \frac{1}{N_v} \sum_{j=1}^{N_v} \delta s_j$.

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.

We can now compare our observations in 2H-NbSe$_2$ with the one-dimensional model discussed above. The barrier to overcome during pinning by thermal fluctuations is given by Arrhenius’ law $U_b = k_B T \ln(\omega t)$. 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 order of $(U_b/k_B)/\ln(\omega t)$. Note that $U_b$ denotes here the energy barrier for vortex creep, in contrast to the pinning energy of one defect site$^{37,38}$. With $U_b \sim k_B T c_c(j_c/j_{dp})/G_i^{1/2}(B/H_s)^{1/2}$, the theory of weak collective pinning$^{24,39}$ provides an estimate for this energy scale. Here $j_{dp} = e\Phi_0/12\sqrt{\pi^2\lambda^2\xi}$ denotes the depairing current and the Ginzburg-Levanyuk number $G_i \sim [T_c/H_s(0)\xi(0)]^{1/2}$ measures the importance of thermal fluctuations with $H_s(0)\xi(0) = \Phi_0^2/(8\pi^2\lambda(0)^2$ the superconducting condensation energy. We estimate $j_c/j_{dp} \sim 10^{-4}$ in 2H-NbSe$_2$ and $G_i \sim 10^{-4}$, from Refs. $^{40-45}$, and obtain $U_b \sim 10 k_B T_c$, which is compatible with $U_b \lesssim k_B T_c \ln(\omega t)$, provided $\omega t \approx 2 \times 10^4$. This is somewhat larger than the values considered above. Given the simplicity of the one-dimensional model, the agreement is still remarkable. All important features predicted by the model (Fig. 2)—the disappearance and reappearance of the directed motion, together with the temperature-evolution of the jitter motion—are found in the experiment (Fig. 4).

Given that the experimental time scale spans several minutes, our observation $\omega t \sim 10^4$ suggest a value for $\omega$ of order of one Hz. While a route for accurate determination of this attempt frequency is still lacking, estimates for a single vortex relate to the Labusch parameter $\alpha_L$ (effective$^{47}$ pinning curvature) and the vortex viscosity $\eta$ via $\omega = \alpha_L/\eta$. While values in the range $10^4-10^5$ Hz have been reported$^{48}$ the analysis assumes vibrations with large $k$-vectors. In our case, vortices are not isolated, but rather interact non-locally with many vortices$^{2,3,49-51}$. Low $k$-vectors, or wave-lengths comparable to the sample size, leads to highly dispersive elastic moduli which modify the attempt frequency by orders of magnitude$^{50,52,53}$. 

Similar to our observation, previous measurements of slow vortex dynamics have reported$^{51,54}$ very low frequency values for thermal motion and creep. Also, creep rates observed in layered cuprate superconductors involve extremely large time scales, indicating the relevance of collective creep$^{2,3,52,54}$. Similar results are found in stochastic behavior of particle arrangements$^{55}$, where, depending on particle interactions, the dynamics transits from individual random motion to flocking. The time scale related to flocking motion, shows a divergent behavior with increasing interaction. In analogy, even if the attempt rate of individual vortices is large, the dynamics as a lattice involves rates that are many orders of mag-

![FIG. 4.](image-url)
nitude smaller. Another implication of this observation is that the temperature is far from melting, thus favoring a collective rather than a single-vortex dynamics. It is this near-equilibrium configuration with ultra-small collective dynamics that allows for the observed cooling imposed creep in our experiments.

In conclusion the model of self-imposed creep explains the critical state dynamics in 2H-NbSe$_2$ 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 many other uniaxial superconductors with weak pinning. Among the different dynamical phenomena of the vortex state, self-imposed creep exemplarily shows the important roles played by history and state preparation. The appearance of creep when cooling provides a new opportunity to consider self-imposed thermal effects in other systems, such as alloys and mixtures like concrete or rocks, liquids re-solidifying under stress, steel under stress, colloidal systems or skyrmions. For all these complex systems vortex matter continues to serve as an ideal testing ground.

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