Pinning-Induced Folding-Unfolding Asymmetry in Adhesive Creases

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The compression of soft elastic matter and biological tissue can lead to creasing, an instability where a surface folds sharply into periodic self-contacts. Intriguingly, the unfolding of the surface upon releasing the strain is usually not perfect: small scars remain that serve as nuclei for creases during repeated compressions. Here we present creasing experiments with sticky polymer surfaces, using confocal microscopy, which resolve the contact line region where folding and unfolding occurs. It is found that surface tension induces a second fold, at the edge of the self-contact, which leads to a singular elastic stress asymmetry that is caused by contact line pinning, in a way that resembles wetting of liquids on imperfect surfaces. Contact line pinning is therefore a key element of creasing: it inhibits complete unfolding and gives soft surfaces a folding memory.

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Creases are ubiquitous to nature and can readily be observed by closing ones’ hand or bending ones’ arm: Soft tissue responds to compression by folding into deep valleys of self-contacting skin [1–3]. The morphology of mammalian brains [4–7] or tumors [8] is dominated by creases that emerge from tissue growth under constraint conditions. Similarly, polymer coatings in technological applications may suffer from creasing due to swelling [9–13], but also compressed elastomers [1,14–16] and viscoelastic liquids [15] display this instability.

The canonical creasing behavior can be realized by uniaxially compressing a slab of a soft material [Fig. 1(a)]. Sharp creases form via a subcritical bifurcation after a reaching a critical strain $\epsilon_c$. When subsequently releasing the strain to below $\epsilon_c$, the length of the crease is only gradually reduced. This implies bistability since either creased or homogeneous states can be found, depending on the deformation history [17–19]. Recent studies clarified the onset of creasing [2,20,21], invoking surface tension [17,18,22] or the presence of a skin [23] to explain the observed bistability.

Interestingly, a microscopic residual crease typically remains even after the strain is fully released to $\epsilon = 0$ [Fig. 1(a)]. This small feature, referred to as a “scar,” is of great significance, since it serves as a nucleus for creases when repeating the compression [11,17,18]. As such, these scars endow soft materials with mechanical memory [11,17], offering a potential of dynamic programmability of surface folds. Despite their importance, scars have remained somewhat enigmatic. It was found that scars are not due to material failure, and their persistence was argued to originate from adhesion [16–18]. However, it is not clear whether adhesion and surface tension can actually lead to a reduction of surface energy compared to the flat, scarless state. Recent work focused on the consequences of surface tension to the onset bifurcation of creasing [22], but it is not known how surface tension affects folding and unfolding at the microscale.

In this Letter we resolve the micro- and macromorphology of adhesive creases by confocal microscopy [Fig. 1(b)], and identify the role of surface tension ($\gamma$), inside the self-contact. It is found that surface tension induces a second fold at the edge of the self contact, turning the surface into a T-shaped profile [Fig. 1(b), bottom right]. This involves a change of the contact angle $\theta$ from 180° for $\gamma \sim 0$ to 90° for $\gamma > 0$. Further, we show that folding and unfolding the crease are, at the microscale, intrinsically asymmetric processes. The unfolding is inhibited, and ultimately prevented, by contact line pinning. This pinning-induced hysteresis implies a new type of bistability, even far above the onset of creasing, and offers a natural explanation for the formation of scars.

Experimental.—A layer of a soft polymer gel (Dow Corning CY52-276, components A:B mixed 1:1 or 1.4:1
Elastocapillary self-contact.—We first quantify the morphology of the elastocapillary self-contact, and how it is altered by surface tension. Figure 1(b) shows creased surface profiles from three experiments with different elastocapillary lengths: $\ell = \gamma / G \approx 0$ (red), $\sim 20 \mu$m (green, $G \approx 1.1$ kPa), and $\sim 300 \mu$m (blue, $G \approx 65$ Pa). Increasing the elastocapillary length amplifies the relative importance of surface tension, which on the macroscale leads to shallower indentations.

However, the most important consequence of $\gamma > 0$ is reflected in the micromorphology of the contact region (high magnification data, insets of Fig. 1(b)). The angle $\theta$ that the surface describes at the contact line changes from a gentle touchdown with $\theta = 180^\circ$ for $\gamma \sim 0$, to $\theta = 90^\circ$ for $\gamma > 0$. Hence, besides the fold at the bottom of the crease, there is a second fold at the top of the self-contact where the surface profile acquires a T shape. On scales much smaller than $\ell$, capillarity will be dominant over elasticity and we expect the contact angle to be given by Neumann’s law [26–28]. In case of perfect self-adhesion, which is expected for soft polymer gels, the surface tension in the self-contact vanishes completely. The remaining balance of gel-liquid surface tensions implies the fold of $\theta = 90^\circ$ [Fig. 1(b)], in good agreement with our experiments.

Consequent to this second fold is a strong curvature $\kappa$ of the free surface near the contact line. When attempting to quantify this curvature, however, it turns out that $\kappa$ does not reach a well-defined limiting value. Rather, it still grows as we reach our measurement resolution, which is well below the elastocapillary length. Figure 2 (purple markers) shows the measured curvature as a function of the distance to the contact line, for $\ell \sim 20 \mu$m. Surprisingly, the data suggest a logarithmic divergence of $\kappa$ as the contact line is approached.

The logarithmic singularity of curvature is caused by the capillary boundary condition, which forces the material into a $90^\circ$ angle. Analogous to the bottom of the crease [2,19,20], a fold of angle $\theta$ introduces a weak (logarithmic) stress singularity. For a neo-Hookean solid the stress singularity reads $\sigma_\theta = (\pi / \theta - \theta / \pi) G \ln(|x|) + p_0$ [29], where $\theta$ is the fold angle in radians and $p_0$ a gauge pressure. At the bottom of the self contact, where the crease was initiated, the fold has an angle $\theta = 2\pi$. In contrast, the two folds at the contact line involve an angle $\theta = \pi / 2$, connecting the self-contact to the free surface in a T shape. Thus the elastic stress must be balanced by the Laplace pressure $\gamma \kappa = -p_\text{el}$, which for a right angle gives the elastocapillary balance.
The curves reflect differences in gauge pressure ($\kappa$) and in simulations (lines). The horizontal shift between the curves reflects differences in gauge pressure $p_0$ in the log, which depends on the ratio $l/L$. The experimental data was taken after compressing, for $l \sim 20 \mu$m.

$$\gamma \kappa \approx \frac{3}{2} G \ln \left( \frac{b}{|x|} \right),$$

valid at distances $|x| \ll l$. Here we absorbed the gauge pressure $p_0$ into a dimensionless constant $b$. The prediction (1) is shown as the solid line in Fig. 2, in excellent agreement with the experimental data. We only adjusted the value of $b$, which cannot be derived by the local analysis of the singularity: it reflects a gauge pressure that is inherited from the full solution, invoking scales larger than $l$.

To explore this effect, we performed finite element simulations of elastocapillary creases, assuming a neo-Hookean solid with constant surface tension (see Supplemental Material for details [24]). Numerical results are obtained for various ratios of $l/L$, where $L$ is the crease length. All numerical results exhibit the logarithmic divergence of curvature (Fig. 2, lines), following the prediction (1). The fact that the data are shifted laterally reflects the nonuniversality of $b$. Still, the numerical data gradually approach the experimental data as $l/L$ approaches the experimental value.

Thus we conclude that the capillary nature of adhesive creases has two important consequences at the microscale: (i) it governs the contact angle $\theta$, implying a second fold at the top of the crease where the surface is $T$ shaped. (ii) the fold, in turn, introduces a logarithmic singularity of the elastic stress and the surface curvature at the edge of the self-contact.

Intermediate asymptote.—Contrarily, for $x \gg l$, we expect capillary effects to play no role and hence, the problem should be described by a purely elastic solid mechanics. Then the crease length $L$ becomes the relevant scale, and the morphology of the free surface was predicted to exhibit the scaling [19]

$$\frac{y - y_0}{L} \sim \left( \frac{x}{L} \right)^{2/3},$$

where $y_0$ is the vertical coordinate of the contact line. This intermediate asymptote is expected to be valid whenever $l \ll x \ll L$. Figure 3 shows the collapse of simulations and measurements for $\gamma \approx 0$ and various $\epsilon$, and confirms the $2/3$ exponent.

Folding-unfolding asymmetry.—Elastocapillary (inner) and elastic (intermediate) regions can be collapsed simultaneously by choosing the appropriate scales $x = l_x \hat{x}$ and $y = l_y \hat{y}$ on each axis. The inner, elastocapillary morphology (1), requires $l_x/l_y = \ell'$ to collapse the data. The scaling behavior of the intermediate region (2), requires $l_y/l_x^{2/3} = L^{1/3}$. Both requirements are fulfilled simultaneously by choosing

$$l_x = \ell'^{3/4} L^{1/4}, \quad l_y = \ell'^{1/2} L^{1/2}.$$
observed during folding are manifestly different from the profiles during unfolding, the latter being more shallow. This asymmetry between folding and unfolding is not observed in the simulations. For a given set of parameters, the numerical minimization of elastocapillary energy selects a unique crease morphology. Clearly, an element beyond equilibrium elastocapillary mechanics is required to properly interpret the experiments, which originates from contact line physics.

Contact line pinning.—The difference in morphologies upon switching from compression to expansion is visualized in the inset of Fig. 5. The green profile is taken during the compression phase, while the red profile is obtained after a subsequent release of strain. Clearly, the global indentation of the free surface decreased upon reducing \( \epsilon \), while the crease length \( L \) remained identical up to the measurement resolution. This offers direct evidence for contact line pinning, where there is no “unfolding” at all, i.e., no change of material points at the contact line [cf. sketch in Fig. 1(a)].

In the present context, the morphological hysteresis readily impacts the crease length \( L \), which is no longer a pure function of the imposed strain \( \epsilon \). This is shown in Fig. 5 (main plot), reporting \( L \) versus \( \epsilon \). While bistability is well known below the onset of creasing, we here find an additional bistability that is caused by contact line pinning: the shaded area shows that, even well above the onset of creasing, the crease length still exhibits a history dependence. After increasing \( \epsilon \), smaller \( L \) are selected, rather than after decreasing \( \epsilon \), especially for \( \ell \sim 20 \) \( \mu \)m (orange curve, same data as in Fig. 4). For \( \gamma \sim 0 \), the measurement error is comparable to the detected hysteresis, and no definitive statement on its existence can be made. Because of experimental limitations we can only give an upper bound for the scar length \( \sim 10 \) \( \mu \)m. However, the fact that this hysteresis loop becomes more pronounced for larger \( \ell \) suggests a capillary origin of contact line pinning.

The observed contact line pinning is analogous to the wetting of liquids on solid surfaces [30]. In that case, the motion of solid-liquid-vapor three-phase contact lines is arrested by pinning on heterogeneities of the solid surface. Such heterogeneities lead to a complicated energy landscape which allows for a range of stable liquid morphologies, leading to a range of contact angles. The pinning observed for creasing could be of similar origin as in wetting, where an energetic barrier \( \Delta \gamma \) is required to move the contact line across features of the surface topography. We therefore expect a range of mechanically stable crease lengths of the order of \( \Delta L \sim \Delta \gamma / G \). This ultimately prevents complete unfolding: the remaining elastic energy is not sufficient to overcome this pinning barrier. This explanation for scars is similar to that given in Ref. [18], which was phrased in terms of an energy release rate rather than \( \Delta \gamma \). From our findings, however, it is clear that the
energy release rate for unfolding an adhesive crease cannot be given by the reversible work of self-adhesion, based on surface energies, as that would not offer a mechanism for contact line pinning. Indeed, our finite element simulations (with reversible adhesion) do not give a folding-unfolding asymmetry. They lack the new bistability indicated by the shaded areas in Fig. 5, and do not admit any scars. To exclude the tracer particles as cause for the scars in our experiments, we also tested an uncoated specimen by brightfield reflection microscopy, finding the same behavior (see Supplemental Material [24]). Our explanation in terms of contact line pinning is also consistent with the observation of scar annealing on long timescales [16,18]: once again, this resembles the case of liquid wetting, where contact line pinning can indeed be overcome by thermal activation [31–33].

Outlook.—Our confocal microscopy experiments have revealed an intrinsic folding-unfolding asymmetry, induced by contact line pinning, that offers a natural explanation for scars. Pinning is therefore a central element in creasing that needs to be accounted for in theory and simulations, and offers a way to articulate the role of “defects.” In addition, we have shown that surface tension dictates the mechanics at scales below the elastocapillary length, folding the solid into well-defined contact angles. As such, the crease morphology opens a new route to quantify solid-solid interaction of stimuli-responsive hydrogels, Nat. Mater. 9, 159 (2010).

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Correction: The omission of a marker indicating “Featured in Physics” has been fixed.