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Nonlinear plastic modes in disordered solids

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We propose a theoretical framework within which a robust micromechanical definition of precursors to plastic instabilities, often termed soft spots, naturally emerges. They are shown to be collective displacements (modes) \( \hat{\gamma} \) that correspond to local minima of a barrier function \( b(\hat{\gamma}) \), which depends solely on inherent structure information.

We demonstrate how some heuristic searches for local minima of \( b(\hat{\gamma}) \) can \textit{a priori} detect the locus and geometry of imminent plastic instabilities with remarkable accuracy, at strains as large as \( \gamma_c - \gamma \sim 10^{-2} \) away from the instability strain \( \gamma_c \). Our findings suggest that the \textit{a priori} detection of the entire field of soft spots can be effectively carried out by a systematic investigation of the landscape of \( b(\hat{\gamma}) \).

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Plastic flow of disordered solids subjected to external loading is known to occur via localized rearrangements of small sets of particles, coined shear transformations [1]. Such rearrangements have been identified in experiments on bubble rafts [2], foams [3], emulsions [4,5], and colloidal glasses [5,6], as well as in atomistic computer simulations of model glasses [7,8]. An example of such a shear transformation, observed in a model glass in two dimensions deformed under athermal quasistatic shear [9], is displayed in Fig. 1(b). Shear transformations are known to self-organize in spatially correlated patterns [10–16] in solids subjected to large stresses and low deformation rates. Their densities and other statistical properties, and mechanical consequences, are a subject of much recent debate [17–26]. Two questions, central to theoretical descriptions of elastoplasticity, that we address in this work are whether shear transformations can be predicted \textit{a priori} and if so, how.

The micromechanical process in which an athermal disordered solid destabilizes under quasistatic deformation is understood asymptotically close to an instability strain \( \gamma_c \) as a saddle-node bifurcation of the potential energy \( U \) [12,27,28]. The immediate precursors to shear transformations at strains \( \gamma \to \gamma_c \) are identified as destabilizing eigenfunctions \( \hat{\Psi} \) (i.e., their associated eigenvalues vanish at \( \gamma_c \)) of the dynamical matrix \( M_{ij} = \frac{\partial^2 U}{\partial x_i \partial x_j} \), where \( \vec{x} \) denotes the coordinate vector of the \( i \)-th particle. Such an eigenfunction is presented in Fig. 1(a) [9]. In the following we refer to such eigenfunctions as destabilizing modes to distinguish them from the postinstability displacements of particles (agglomerations of shear transformations) that can be spatially extended. In Fig. 1(b) we demonstrate that, when the postinstability displacements are not spatially extended, but rather form an isolated elementary shear transformation, their spatial structure is very similar to that of the destabilizing mode. In contrast with the postinstability displacements that depend in general on a specific choice of dynamics [14] and on external control parameters such as temperature [15], strain [11], and strain rate [16], the spatial structure of destabilizing modes is an intrinsic characteristic of the multidimensional potential energy function and is therefore the focus of the present study.

FIG. 1. Plastic instability in a sheared two-dimensional model glass. (a) Destabilizing eigenfunction \( \hat{\Psi} \). (b) Elementary shear transformation: the postinstability displacements that followed the instability of (a). (c) Nonaffine displacements \( \vec{v} \) calculated at \( \delta \gamma \equiv \gamma_c - \gamma = 0.007 \). This delocalized field is used as the initial conditions \( \vec{z}_{\text{init}} \) for the minimization of \( b(\hat{\gamma}) \) (see the main text), the result of which is the plastic mode \( \vec{\xi} \) displayed in (f). (d) and (e) Intermediate states along the minimization of \( b(\hat{\gamma}) \).

A robust mechanical definition of the precursors of plastic instabilities away from instability strains has not yet been put forward. Much effort has been dedicated recently to studying...
the role played by low-frequency normal modes in determining these precursors [29–34]. One key difficulty encountered in such studies is that low-frequency plane waves, which have no appreciable effect on plasticity [17], dominate the lower parts of the spectra of conventional model glasses, thus hindering attempts to use low-frequency modes to define flow-defect densities and correlate them with rates of plastic flow.

Another difficulty, which has been largely overlooked in the context of elastoplasticity, is that mere frequencies of normal modes are not indicative of their relevance to plastic processes. Here we show that the effective detection of such modes away from plastic instabilities is possible by accounting for the relevant nonlinearities of the potential energy landscape. We provide a theoretical framework that naturally embeds a micromechanical definition of the precursors to plastic instabilities and that effectively accounts for the said nonlinearities.

We begin the discussion by considering an athermal elastic solid, of $N$ particles in $d$ dimensions, and let $\hat{z}$ denote an $Nd$-dimensional unit vector, i.e., $\hat{z}_i : \hat{z}_j = 1$. Here and in what follows repeated indices, labeling particles, are understood to be summed over unless indicated otherwise. The coordinates $\vec{r}$ are displaced in the direction defined by $\hat{z}$ according to $\delta \vec{r} = s \hat{z}$ and we expand the potential energy $U$ as

$$\delta U_{\hat{z}}(s) = U(s) - U_0 \approx \frac{1}{2} \kappa_s s^2 + \frac{1}{6} \tau_s s^3,$$

where $U_0$ is the energy of the minimum in which the system resides, $\kappa_s \equiv M_{ij} : \hat{z}_i \hat{z}_j$ is the stiffness associated with $\hat{z}$, and $\tau_s = \partial^2 U / \partial \vec{r}_s \partial \vec{r}_s \hat{z}_s \hat{z}_s \hat{z}_s$ is referred to in the following as the asymmetry associated with $\hat{z}$. Within this cubic expansion, stationary points occur at $s = 0$ and $s = -\frac{1}{3} \tau_s^{-1}; s = 0$ corresponds to the minimum in which the system resides, while $s_0$ represents the saddle point (energy barrier) that separates this minimum and a neighboring inherent state. We thus define the energy difference between these stationary points, within the cubic expansion, as our barrier function

$$b(\hat{z}) \equiv \frac{1}{2} \kappa_s s_0^2 + \frac{1}{6} \tau_s s_0^3 = \frac{2\kappa_s}{3\tau_s}. \tag{2}$$

We emphasize that $b(\hat{z})$ is defined for a particular inherent state of an elastic solid in mechanical equilibrium and is a function of the multidimensional direction $\hat{z}$. It has a rough landscape in this work we focus on directions $\hat{\pi}$ that correspond to local minima of $b(\hat{z})$, i.e., they satisfy $\partial b / \partial \vec{r} \big|_{\vec{r}=0} = 0$, and $\partial^2 b / \partial \vec{r} \partial \vec{r} \big|_{\vec{r}=0}$ is positive semidefinite. We refer to these directions in what follows as plastic modes. From the definition of $b(\hat{z})$ it is clear that plastic modes $\hat{\pi}$ are associated with small stiffnesses $\kappa_{\pi}$ and large asymmetries $\tau_{\pi}$; they can be found numerically by minimizing $b(\hat{z})$ over directions $\hat{z}$, starting from some initial direction $\hat{z}_{\text{ini}}$, as demonstrated in Figs. 1(c)–1(f). Small $b(\hat{z})$’s should appropriately describe low saddle points (barriers) that separate the system from neighboring inherent states. We therefore expect modes $\hat{\pi}$ that correspond to low-lying minima of $b(\hat{z})$ (which are found by cleverly choosing an appropriate $\hat{z}_{\text{ini}}$ for the minimization) to encode information about imminent plastic instabilities.

This approach is demonstrated in Fig. 1: in Fig. 1(a) we display a destabilizing mode $\hat{\Psi}_1$, calculated at the first-encountered plastic instability in an athermally sheared model glass, here at a strain $\gamma_c = 0.011521$. To this instability, at strains $\gamma = \gamma_c - \delta \gamma$, the nonaffine displacement field $\vec{v}_1 \equiv -M_{ij}^{-1} \cdot \partial^2 U / \partial \vec{r}_i \partial \vec{r}_j \big|_{\vec{r}=0}$ is calculated [27]. An example of $\vec{v}$, calculated at $\delta \gamma = 0.0007$, is shown in Fig. 1(c). At this distance (in strain) from the instability, the nonaffine displacements $\vec{v}$ are largely delocalized. We use the normalized $\hat{v} = \vec{v} / \|\vec{v}\|$ as the initial conditions $\hat{z}_{\text{ini}}$ for the minimization of $b(\hat{z})$; snapshots along the minimization are displayed in Figs. 1(d) and 1(e).

Under convergence, we find a local minimum in the direction $\hat{\pi}$, which is displayed in Fig. 1(f). The resemblance between $\hat{\pi}$ and the destabilizing mode $\hat{\Psi}_1$ is striking: Both the geometry and the core location appear to agree perfectly. This remarkable agreement is quantified in the inset of Fig. 2(b), where we plot the overlaps $\hat{\Psi}_1 \cdot \hat{\pi}$ as a function of the distance to the instability strain.

The protocol described above is carried out over a broad range of intervals $\delta \gamma$, as specified in the caption of Fig. 2. For each $\delta \gamma$, after finding $\hat{\pi}$ as described above, we calculated its associated energy variation $\delta U_{\hat{\pi}}(s)$ and stiffness $\kappa_{\pi}$, which are displayed in Figs. 2(a) and 2(b), respectively. In this example, already at a distance of the order $\delta \gamma \approx 10^{-3}$ to the instability strain, following the plastic mode $\hat{\pi}$ would carry the system above an energy barrier and into a neighboring minimum.

The resolution of the plastic mode as shown in Fig. 1 uses the nonaffine displacements $\vec{v}$ as the heuristic guess for $\hat{z}_{\text{ini}}$: this choice is made to demonstrate the usefulness of the framework: Despite the extended character of $\vec{v}$, it has a large overlap with the plastic mode $\hat{\pi}$ and thus resides in the basin of $\hat{\pi}$ on the landscape of $b(\hat{z})$. Obtaining the full field of plastic modes, however, requires using other heuristic $\hat{z}_{\text{ini}}$’s, which reside in basins that belong to other plastic modes. We leave the investigation of the optimal heuristics for the detection of the full field of plastic modes for future work.
We also plot in Fig. 2(b) the stiffness $k_{\bar{\psi}}$ associated with the destabilizing mode $\bar{\psi}_\tau$. We find that only very close to the instability ($\delta \gamma \lesssim 10^{-5}$), the scaling $k_{\bar{\psi}} \sim \sqrt{\delta \gamma}$ holds, whereas the stiffness associated with $\bar{\pi}$ follows $k_{\bar{\pi}} \sim \sqrt{\delta \gamma}$ up to strains of order 1% away from the instability. This finding supports the robustness of our definition of plastic modes and the usefulness of our framework. It also supports the picture proposed by a number of recent studies [19–24] that assumes that the (reversible) destabilization process of a soft spot in a deformed glass is predominantly coupled to the external load and not to other coexisting (reversible) destabilization processes.

The scaling $k_{\bar{\pi}} \sim \sqrt{\delta \gamma}$ can be derived as follows. Modes $\bar{\pi}$ pertain to local minima of $b(\hat{\pi})$ and therefore satisfy $\left. \frac{\partial b}{\partial \pi} \right|_{\pi=\hat{\pi}} = 0$, which implies that (see [9])

$$\frac{\partial^2 U}{\partial \pi_j \partial \pi_k} : \bar{\pi}_j \bar{\pi}_k = \frac{\tau_\pi}{k_{\bar{\pi}}} M_{ij} \cdot \bar{\pi}_j.$$

Using this relation, we calculate the leading-order variation of the stiffness $k_{\bar{\pi}}$ with strain as

$$\frac{dk_{\bar{\pi}}}{d\gamma} \simeq \frac{dM_{ij}}{d\gamma} : \bar{\pi}_j \bar{\pi}_k \simeq \frac{\partial^2 U}{\partial \pi_j \partial \pi_k} : \bar{\pi}_j \bar{\pi}_k = \frac{\tau_\pi}{k_{\bar{\pi}}} M_{ij} \cdot \frac{\partial^2 U}{\partial \pi_j \partial \gamma}.$$

As $\gamma \to \gamma_c$, $k_{\bar{\pi}} \to 0$, but $\tau_\pi \sim \sqrt{\delta \gamma}$ goes to a constant, yielding the differential scaling relation $\frac{dM_{ij}}{d\gamma} \sim -\frac{1}{\kappa_{\bar{\pi}}}$ and thus the observed scaling $k_{\bar{\pi}} \sim \sqrt{\delta \gamma}$.

Comparison to normal modes. How indicative are normal modes of immanent plastic instabilities, compared to plastic modes? In Fig. 3(a) we present a scatter plot of the barrier function evaluated at normal modes $\bar{\psi}_\tau$ vs their square of their associated frequencies $\omega^2$, calculated for a few tens of undeformed (isotropic) solid realizations. A clear trend appears: Smaller values of $b(\bar{\psi}_\tau)$ are found for lower-frequency modes. The circled data point represents the mode $\bar{\psi}_{\text{min}}$ associated with the lowest value of $b(\bar{\psi}_\tau)$ among all modes calculated; it is displayed in Fig. 3(d). Remarkably, this normal mode displays the same spatial features as observed for destabilizing modes, reinforcing that $b(\bar{\pi})$ is indeed sensitive to plasticlike modes. The variation $\delta b(\bar{\psi}_{\text{min}})(s)$ is plotted in Fig. 3(b) (solid line). Despite possessing the smallest $b$ among our entire ensemble of modes, $\delta b(\bar{\psi}_{\text{min}})(s)$ displays only a slight asymmetry between positive and negative displacements $s$ and the energy monotonically increases with $|s|$. Using $\bar{\psi}_{\text{min}}$ as the initial condition $\bar{\psi}_{\text{init}}$ for the minimization of $b(\bar{\psi})$, we find the plastic mode $\bar{\pi}$ displayed in Fig. 3(e). On the face of it, $\bar{\psi}_{\text{min}}$ and $\bar{\pi}$ appear to be very similar in their spatial structure and geometry. However, examining the corresponding variation $\delta U_{\bar{\pi}}(s)$, represented by the dashed line in Fig. 3(b), reveals a dramatic difference between them: Following $\bar{\pi}$ takes the system over a energy barrier to a neighboring minimum.

We further utilize our ensemble of normal modes to study the relation between the degree of localization of modes and their associated asymmetries $\tau_{\bar{\psi}}$. A similar analysis was carried out in [33] in the context of the unjamming point [35–37]. We quantify the degree of localization of a mode $\bar{\psi}$ via its participation ratio $e = \left| \sum_{i=1}^{N} (\bar{\psi} \cdot \bar{\psi}_i)^2 \right|^{-1}$; localized modes have $e \sim N^{-1}$, whereas maximally delocalized modes have $e \sim 1$. In Fig. 3(c) we plot the means $\left| \tau_{\bar{\psi}} \right|$, averaged over modes $\bar{\psi}$ with similar participation ratios, for systems of $N = 1024$ and 4096. We find that for participation ratios $e < 10^{-1}$, the asymmetries follow $\left| \tau_{\bar{\psi}} \right| \sim (eN)^{-1}$. This can be explained with the following simple model: If there are effectively $N^a$ nonzero components in a normal mode ($0 \leq a < 1$), normalization then requires that a characteristic nonzero component is of magnitude $\left| \bar{\psi}_i \right| \sim N^{-a/2}$. The participation ratio is then expected to follow $e \sim N^{a-1}$ (due to summing over positive terms). Since the pairwise potential is short ranged and the tensor elements $U_{\bar{\pi},\bar{\pi}_i}$ are of either sign, then $\tau_{\bar{\psi}}$ consist of a sum over $N^a$ terms, each of order $\left| \bar{\psi}_i \right| \sim N^{-3a/2}$, of random signs and we therefore expect $\left| \tau_{\bar{\psi}} \right| \sim N^{a-1}$, in consistency with our measurement. For participation ratios $e > 10^{-1}$, this relation breaks down and asymmetries are much smaller than what is predicted by this simple model, which assumes that normal modes are random objects. Nevertheless, the same trend remains unchanged: Delocalized modes are associated, on average, with more

In a system of linear size $L$, the scaling law $k_{\bar{\pi}} \sim \sqrt{\delta \gamma}$ is only expected to hold below $\delta \gamma \sim L^{-4}$ due to hybridizations with low-frequency plane waves.

3We consider the absolute magnitudes $|\tau_{\bar{\psi}}|$ since normal modes $\bar{\psi}$, and therefore also their associated asymmetries $\tau_{\bar{\psi}}$, are defined up to a sign.
Figs. 4(c) and 4(d) two additional examples of plastic modes depends on the loading conditions imposed on the solid. In to which they correspond. 

Structural features that do not depend on the particular minima associated with different local minima of \( \Psi \). Random \( \hat{\psi} \) modes found by choosing a random \( \hat{\psi} \) that correspond to local minima of the barrier function \( b(\hat{\psi}) \), coined plastic modes, are indicative of directions in configuration space that lead to plastic instabilities and more so compared to the most localized low-frequency normal modes. As such, our framework can serve as a solid basis for instability-detection algorithms. Such algorithms are highly desirable, as they can put to test theoretical frameworks of elastoplasticity that involve the dynamics of a population of soft spots. These algorithms need not be restricted to the investigation of plastic flow in disordered solids; the generality of our framework would render them suitable for studying a diverse set of systems, including dislocated crystalline solids, deeply supercooled liquids, and proteins.

Furthermore, our theoretical framework explains the origin of the localized nature of plastic instabilities. Building on our framework, we predict that the stiffness associated with plastic modes follows \( \kappa \sim \sqrt{\gamma_c - \gamma} \) and show numerically that this scaling holds over a large range of strains away from an instability strain \( \gamma_c \). This adds relevance to recently proposed models that assume that reversible destabilization processes of soft spots are decoupled from each other. Finally, we have investigated the spatial features of plastic modes and provided evidence that the detailed geometry of plastic modes is sensitive to the loading conditions imposed on the solid.

Our approach demonstrates the usefulness of the concept of exploring the direction space associated with an inherent state of a solid as a means of extracting micromechanical information that is highly relevant to nonlinear flow processes. Similar approaches could likely be applied towards studying mechanical instabilities in, e.g., granular solids \( \gamma \) and towards studying other classes of low-energy excitations in glassy solids, e.g., two-level systems \( \gamma \). We thank Eran Bouchbinder, Smarajit Karmakar, Gustavo Düring, Itamar Procaccia, and Matthieu Wyart for fruitful discussions. This research was funded by the Amsterdam Academic Alliance fellowship of the University of Amsterdam and the Vrije Universiteit Amsterdam.