Comment on "Moving Glass Phase of Driven Lattices"

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Comment on “Moving Glass Phase of Driven Lattices”

In a recent Letter [1] Giamarchi and Le Doussal (GL) showed that when a periodic lattice is rapidly driven through a quenched random potential, the effect of disorder persists on large length scales, resulting in a Moving Bragg Glass (MBG) phase. The MBG was characterized by a finite transverse critical current and an array of static elastic channels.

They use a continuum displacement field \( \mathbf{u}(r, t) \), whose motion (neglecting thermal fluctuations) in the laboratory frame obeys 

\[
\eta \partial_t u_\alpha + \eta \mathbf{v} \cdot \nabla u_\alpha = c_{11} \partial_{\alpha} \nabla \cdot \mathbf{u} + c_{66} \nabla^2 u_\alpha + F^p_\alpha + F_\alpha - \eta v_\alpha, \]

where \( F_\alpha \) is the external driving force. As in [1], we choose \( F_\alpha = F_\alpha^{stat} + F_\alpha^{dyn} \) and denote by \( y \) the \( d-1 \) transverse directions. GL observe that the pinning force \( F_\alpha^{dyn} \) splits into static and dynamic parts, \( F_\alpha^{stat} = F_\alpha^{stat} + F_\alpha^{dyn} \), with 

\[
F_\alpha^{stat}(r, \mathbf{u}) = \rho_{V}(r) \sum_{K \cdot \mathbf{v} = 0} |K_{\alpha} e^{iK \cdot (r - \mathbf{u})} - \rho_{V}(r) \nabla_{\mathbf{v}} V(r),
\]

and 

\[
F_\alpha^{dyn}(r, \mathbf{u}, t) = \rho_{V}(r) \sum_{K \cdot \mathbf{v} \neq 0} |K_{\alpha} e^{iK \cdot (r - \mathbf{v} - \mathbf{u})} |.\]

GL argue that in the sliding state at sufficiently large velocity \( \mathbf{F}^{stat} \) gives the most important contribution to the roughness of the phonon field \( \mathbf{u} \), with only small corrections coming from \( \mathbf{F}^{dyn} \). Since \( \mathbf{F}^{stat} \) is along \( y \) and only depends on \( u_y \), they assume \( u_x = 0 \) and obtain a decoupled equation for the transverse displacement \( u_y \). Analysis of this equation then predicts the moving glass phase with the aforementioned properties.

In this Comment, we show that the model of Ref. [1] neglects important fluctuations that can destroy the periodicity in the direction of motion. Following recent work by Chen et al. [2] for driven charge density waves, it can be shown [3] that the longitudinal dynamic force \( F_{z}^{dyn} \) does not average to zero in a coarse-grained model, but generates an effective random static drag force \( f_x(r) \). This arises physically from spatial variations in the impurity density, and can be obtained by using a variant of the high-velocity expansion or by coarse-graining methods. To leading order in \( \frac{1}{F} \) its correlations are 

\[
\langle f_x(r) f_x(0) \rangle = \Delta \delta(\mathbf{r}), \quad \Delta = \Delta_z \sim \Delta^2/F, \quad \delta(\mathbf{r}) = \text{variance of the quenched random potential} V(r).\]

The crucial difference from Ref. [1] is that in contrast to \( \mathbf{F}^{stat} \), the effective static drag force \( f_x(r) \) is strictly \( u \)-independent, as guaranteed by the precise time-translational invariance of the system coarse-grained on the time scale \( \sim 1/v \).

In the presence of \( f_x \), we now reexamine both the elasticity and the relevance of longitudinal dislocations (i.e. those with Burger’s vectors along \( x \)). An improved elastic description begins with the equation

\[
\eta \partial_t u_\alpha + \eta \mathbf{v} \cdot \nabla u_\alpha = c_{11} \partial_{\alpha} \nabla \cdot \mathbf{u} + c_{66} \nabla^2 u_\alpha + \delta_{\alpha x} F_x^{stat}(u_y) + \delta_{\alpha x} f_x(r) \tag{1}
\]

Surprisingly, a simple calculation leads to a transverse correlator, \( B_y(r) = \langle [u_y(r) - u_y(0)]^2 \rangle > 0 \), that is (for \( d > 1 \)) asymptotically identical to that found by GL, which exhibits highly anisotropic logarithmic scaling for \( d \leq 3 \). In contrast, the \( u_x \) roughness is dominated by \( f_x \), and \( B_x(r) = \langle [u_x(r) - u_x(0)]^2 \rangle > 0 \), grows algebraically, \( \sim (\Delta_d/c_{66}) v^{d-3} \), for \( d < 4 \) and \( x < c_{66}/v \), crossing over for \( x > c_{66}/v \) (and \( d < 3 \)) to \( B_x(r) \sim (\Delta_d/c_{66} v^2) v^{3-d} H(c_{66}x/vy^2) \), with \( H(0) = \text{const.} \) and \( H(\varepsilon > 1) \sim \varepsilon^{(3-d)/2} \). We stress that because of \( u \)-dependence of \( f_x \) this power-law scaling for \( B_x(r) \) holds out to arbitrary length scales, in contrast to that for \( B_y(r) \) valid only in the Larkin regime as lucidly discussed by GL [1]. Thus, even within the elastic description, transversal correlations along \( x \) are short-ranged (stretched exponential). Stability with respect to dislocations is more delicate. Nevertheless, arguments analogous to those of Ref. [1] suggest that dislocation unbinding will occur for \( d < 3 \), converting the longitudinal spatial correlations to the pure exponential (liquid-like) form. We stress that this situation corresponds not to \( u_x = 0 \), as assumed in Ref. [1], but rather \( \langle u_x^2 \rangle = \infty \) (indeed, \( u_x \) is multivalued).

We therefore argue that for intermediate velocities (for \( d < 3 \)) a moving vortex solid is organized into a stack of liquid channels, i.e. it is a moving smectic. This is in agreement with structure functions and real-space images from recent simulations [4]. The model for this nonequilibrium smectic state will be the subject of a future publication [4]. An interesting possibility is that at very large velocities, nonequilibrium KPZ type nonlinearities (as in Ref. [2]) might lead to a further transition to a more longitudinally ordered state, with rather different underlying physics from the MBG.

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