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Comment on "Moving Glass Phase of Driven Lattices"

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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 $\mathbf{u}(\mathbf{r},t)$, 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_α is the external driving force. As in [1], we choose $F_\alpha = F\delta_{\alpha,x}$ and denote by y the $d-1$ transverse directions. GL observe that the pinning force F^p_α splits into *static* and *dynamic* parts, $F^p_\alpha = F^{stat}_\alpha + F^{dyn}_\alpha$, with $F^{stat}_\alpha(\mathbf{r},\mathbf{u}) = \rho_0 V(r)\sum_{\mathbf{K}\cdot\mathbf{v}=0} iK_\alpha e^{i\mathbf{K}\cdot(\mathbf{r}-\mathbf{u})} - \rho_0\nabla_\alpha V(r)$ and $F^{dyn}_\alpha(\mathbf{r},\mathbf{u},t) = \rho_0 V(r)\sum_{\mathbf{K}\cdot\mathbf{v}\neq 0} iK_\alpha e^{i\mathbf{K}\cdot(\mathbf{r}-\mathbf{v}t-\mathbf{u})}$. 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^{dyn}_x does *not* average to zero in a coarse-grained model, but generates an effective random static drag force $f_d(\mathbf{r})$. 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 $< f_d(\mathbf{r})f_d(\mathbf{0}) >= \Delta_d\delta(\mathbf{r})$, where $\Delta_d \sim \Delta^2/F$, and Δ is the variance of the quenched random potential $V(\mathbf{r})$. The crucial difference from Ref. [1] is that in contrast to $\mathbf{F}^{dyn}$, the effective static drag force $f_d(\mathbf{r})$ is strictly $\mathbf{u}$-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_d, 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

$$\eta\partial_t u_\alpha + \eta\mathbf{v}\cdot\nabla u_\alpha = c_{11}\partial_\alpha\nabla\cdot\mathbf{u} + c_{66}\nabla^2 u_\alpha + \delta_{\alpha y}F^{stat}_y(u_y) + \delta_{\alpha x}f_d(\mathbf{r}) \qquad (1)$$

Surprisingly, a simple calculation leads to a transverse correlator, $B_y(\mathbf{r}) =< [u_y(\mathbf{r}) - u_y(\mathbf{0})]^2 >$, 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_d, and $B_x(\mathbf{r}) =< [u_x(\mathbf{r}) - u_x(\mathbf{0})]^2 >$ grows algebraically, $\sim (\Delta_d/c^2_{66})r^{4-d}$ for $d < 4$ and $x < c_{66}/\eta v$, crossing over for $x > c_{66}/\eta v$ (and $d < 3$) to $B_x(\mathbf{r}) \sim (\Delta_d/c_{66}\eta v)y^{3-d}H(c_{66}x/\eta v y^2)$, with $H(0) =$ const. and $H(z >> 1) \sim z^{(3-d)/2}$. We stress that because of $\mathbf{u}$-independence of f_d this power-law scaling for $B_x(\mathbf{r})$ holds out to arbitrary length scales, in contrast to that for $B_y(\mathbf{r})$ valid only in the Larkin regime as lucidly discussed by GL. [1] Thus, even within the elastic description, translational correlations along x are short-ranged (stretched exponential). Stability with respect to dislocations is more delicate. Nevertheless, arguments analogous to those of Ref. [4] suggest that dislocation unbinding will occur for $d \leq 3$, 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 $d \leq 3$) 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 [5]. The model for this nonequilibrium smectic state will be the subject of a future publication [3]. 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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