Transition to strictly solitary motion in the Burridge-Knopoff model of multicontact friction

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We show that, in the continuous 1D Burridge-Knopoff model of multicontact friction, motion occurs via stick-slip sliding on a finite length rather than in avalanches, excluding the occurrence of self-organized criticality. We present strong numerical evidence that a transition from collective to strictly solitary motion occurs at a critical value of the interblock interactions. The solitary motion corresponds to successive stick-slip motion of one block between immobile neighbors, repeated periodically in time. This state persists also with open boundary conditions and moderate temperature.

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I. INTRODUCTION

Solid on solid sliding friction is often modeled by one-dimensional spring-block models, meant to represent very different situations. At the atomic scale, friction is well described by the Frenkel-Kontorova, or by the Frenkel-Kontorova-Tomlinson model [1], where the blocks represent individual atoms in interaction with a surface represented as a rigid periodic modulation. At much larger length scales, the Burridge-Knopoff (BK) model, illustrated in Fig. 1, is used to describe sliding tectonic plates. In the BK model, the interaction with the underlying surface is replaced by a phenomenological velocity dependent friction force with a static and a dynamic contribution. The dynamics of tectonic sliding is usually studied by assuming a dynamic friction force that weakens as a function of velocity [2].

In all these models where energy is slowly fed to the system by the moving plate, the dynamics is not uniform but dominated by fast dissipative events corresponding to stick-slip motion of the individual blocks. In their velocity weakening BK model, Carlson and Langer have shown [2] that avalanches of all sizes occur, with a power law size distribution compatible with the empirical Gutenberg-Richter law. This lack of an intrinsic length scale puts this deterministic continuous model into the larger class of systems which are described by the Frenkel-Kontorova, or by the Frenkel-Kontorova-Tomlinson model [1], where the blocks represent individual atoms in interaction with a surface represented as a rigid periodic modulation. Since the finding of Carlson and Langer, the BK model has been studied intensively in this context, particularly in the two-dimensional discretized version proposed by Olami, Feder, and Christensen [5] (OFC). However, several authors claim or suggest that the model does not display self-organized criticality [6–8]. It has even been conjectured that the asymptotic avalanche size distribution is dominated by avalanches of size one, the fraction of larger avalanches converging towards zero as the system size increases [9].

Here we study the multicontact friction variant of the BK model, proposed by Persson [10] to model macroscopic sliding systems in the boundary lubrication regime. The BK model of multicontact friction uses a viscous dynamic friction proportional to velocity, which, contrary to the velocity weakening earthquake models, effectively reduces the range of interactions of the blocks. This approach is justified by previous studies of the same author [11] showing that, at low velocity, a thin lubricant layer exhibits a distribution of pinned solid islands that fluidify and begin to slide when the applied force exceeds a threshold value and pin again as their velocity vanishes.

We find that, after an initial transient, the motion occurs via successive domino-like slipping events of limited size rather than in avalanches, thereby excluding the occurrence of SOC. At a critical value of the interblock interactions close to realistic values for sliding surfaces in the boundary lubrication regime [10], the system reaches a dynamic regime, that we call a solitary state, where the motion occurs via periodic step-like slipping events of single blocks. Surprisingly, the solitary state is not destroyed by open boundary conditions, contrary to the behavior of OFC models [12]. Also the solitary state is robust against small thermal fluctuations.

II. BK MODEL

The BK model of Fig. 1 consists of $N$ blocks of mass $m$ connected, at fixed distances $D$, to a plate moving at constant velocity $v_{f}$ by springs of spring constant $k_{1}$, and to nearest-neighbor blocks by springs of spring constant $k_{2}$ and rest lengths $D$. The plate coordinate is $x=v_{f}t$, and $q_{i}$ is the position of block $i$ with respect to its initial equilibrium position $q_{i}(0)=0$. The force on a block at rest (i.e., $\dot{q}_{i}=0$) is

$$F_{i} = k_{1}(x-q_{i}) + k_{2}(q_{i+1} + q_{i-1} - 2q_{i}).$$

This force is balanced up to a threshold value $F_{th}$ by the static friction force, so that a block remains motionless until it

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{BK_model.png}
\caption{Burridge-Knopoff (BK) model.}
\end{figure}
experiences a force $F_s \geq F_s$. Once in motion, a block is subjected to a viscous force $-2m\gamma \dot{q}_i$. If the block velocity $\dot{q}_i$ vanishes, the static friction force is reintroduced by setting the block velocity to zero if it changes sign. For this reason we always remain in the underdamped regime. The discontinuity of the friction force at $\dot{q}_i=0$ makes the system extremely nonlinear.

We introduce a dimensionless quantity characterizing the dynamic state of block $i$:

$$ h_i = \begin{cases} 0 & \text{if } \dot{q}_i = 0 \text{ (stick)} \\ 1 & \text{otherwise (slip)} \end{cases}. \quad (2) $$

We also introduce

$$ H_i(t) = h_i(t)[h_{i+1}(t) + h_{i-1}(t)] \quad (3) $$
as the number $(0, 1, 2)$ of neighbors slipping while block $i$ is slipping. Note that $H_i=0$ either when block $i$ is at rest ($h_i=0$) or when block $i$ is moving while both neighbors are at rest ($h_{i\pm1}=0, h_i=1$). Since the fraction of time a block is in motion can be quite small, it is useful to average Eq. (3) over a time $\tau$ around $t$

$$ \langle H_i(t) \rangle_{\tau} = \frac{1}{\tau} \int_{t-\tau}^{t} H_i(t')dt', \quad (4) $$
yielding a continuous function, ranging between 0 and 2. By defining $\tilde{h}(t)$ as the fraction of blocks moving at time $t$, the average over all moving blocks

$$ \langle H(t) \rangle_{\tau} = \frac{1}{N} \sum_{i=1}^{N} \langle H_i(t) \rangle_{\tau}, \quad \tilde{h}(t) = \frac{1}{N} \sum_{i=1}^{N} h_i \neq 0 $$

constitutes an order parameter denoting if a system is either in solitary motion, i.e., $\langle H(t) \rangle_{\tau}=0$, or in collective motion, $0 < \langle H(t) \rangle_{\tau} < 2$.

The equations of motion are

$$ m\ddot{q}_i = h_i[-2m\gamma \dot{q}_i + k_1(x - q_i) + k_2(q_{i-1} + q_{i+1} - 2q_i)], $$

where

$$ h_i(t + dt) = \begin{cases} 0, & \dot{q}_i(t)\dot{q}_i(t + dt) < 0 \\ 1, & F_s(t + dt) \geq F_s(t) \end{cases}, \quad (6) $$

$$ h_i(t), \quad \text{otherwise} $$
with $dt$ the time step of numerical integration. The equations of motion are made dimensionless by scaling time by $\sqrt{mk_1}$, positions by $F_s/k_1$ and forces by $F_s$:

$$ \ddot{q}_i = h_i[-2\tilde{\gamma} \dot{q}_i - \tilde{\omega}_0^2 \dot{q}_i + \tilde{k}_2(q_{i-1} + q_{i+1} - 2q_i) + x] \equiv h_i\sigma_i, \quad (7) $$

with $\tilde{\omega}_0 = \sqrt{1 + 2k_2}$ and $\sigma_i$ denoting the total force on block $i$ irrespective of its dynamic state $h_i$. Note that $\tilde{k}_2$ and $\tilde{\gamma}$ are in units of $k_1$ and $\sqrt{k_1/m}$, respectively, and that $\tilde{F}_s=1$. We will only consider dimensionless quantities, and will omit the tilde from now on.

The Eqs. of motion (7) are integrated by a fourth order Runge-Kutta algorithm with time step $dt=0.005$. The initial positions $q_i(0)$ are chosen from a uniform random distribution $q_\approx [-0.005,0.005]$; furthermore $x(0)=0$ and $\dot{q}_i(0)=0$.

We use periodic boundary conditions, unless specified otherwise. The width of the random distribution determines the duration of the transient collective stick-slip behavior. We consider a driving velocity $v_s=0.005$, which is low enough to be in the limit of the transient collective stick-slip behavior. We reproduce Fig. 4 of Ref. [10] extended to larger time. Notice in (c) the transition around $t \sim 20 000$ to solitary motion, causing $\tilde{\sigma}$ and $\tilde{h}$ [see insets of panels (a) and (b)] to become periodic in time. Note that $\Delta t \gg \tau$ in this figure.

\section{III. SOLITARY VERSUS COLLECTIVE MOTION}

In Figs. 2(a) and 2(b) we show the average force $\langle \tilde{\sigma} \rangle$ and the fraction $\tilde{h}$ of moving blocks as in Ref. [10] on a much longer timescale. The initial collective stick-slip behavior is due to the very narrow distribution of forces below $F_s$ at $t=0$. At the first such collective event almost all blocks slip at the same time ($\tilde{h} \approx 1$). As time progresses the distribution of forces $P(\sigma)$ widens and the number of blocks slipping at the same time decreases. After $t \sim 1000$, at any time a number of blocks is moving and, at $t \sim 1800$, the system is said to be in a steady state in Ref. [10].

In the steady state however, the fraction $\tilde{h}$ of moving blocks keeps decreasing, indicating that the system is still equilibrating towards a more favorable state. Finally, at $t = 20 000$, we find that $\tilde{\sigma}$ and $\tilde{h}$ become periodic in time. It is shown in Fig. 2(c) that $\langle H(t) \rangle_{\tau}=0$ when the system becomes periodic. This indicates that blocks slip in a step-like fashion between immobile nearest neighbors ($h_{i\pm1}=0$), whence the name of solitary motion. Once this is the case for the motion of all blocks for longer than the interval between successive slips of the same block, the system is trapped in this solitary state and becomes periodic.

Analytical results give a rationale for this behavior. For solitary motion ($h_{i\pm1}=0$ when $h_i=1$) the equations of motion
(7) become decoupled, and the motion of a single block is that of a discontinuously driven, damped harmonic oscillator. For initial conditions \(q_i(0) = q_{i+1}(0) = 0\) and \(F_i(0) = F_s\) [i.e., \(k_2(q_{i+1} + q_{i-1}) + x = F_s\)], and by assuming \(v_s = \max(q_i)\):

\[
\dot{q}_i + 2\gamma q_i + \omega_0^2 q_i = F_s.
\]

The solution of Eq. (8) for the underdamped case (\(\gamma < \omega_0\))

\[
q_i(t) = \frac{F_s}{\omega_0^2} \left[ 1 - \exp(-\gamma t) \frac{\gamma}{\omega_0} \sin(\omega_0 t) + \cos(\omega_0 t) \right]
\]

reaches zero velocity after a time

\[
\delta t = \frac{\pi}{\omega_0}, \quad \text{with } \omega = \sqrt{\omega_0^2 - \gamma^2}.
\]

In a time \(\delta t\) the block travels a distance [13]

\[
\Delta q = \frac{F_s}{\omega_0^2} \left[ 1 + \exp(-\gamma t/\omega_0) \right].
\]

The interval \(\Delta t\) between consecutive slip events of the same block is given by

\[
\Delta t = \Delta q/v_s,
\]

because, although most of the time a block is not moving, its average velocity has to be equal to the plate velocity \(v_s\). The fraction of time a block is moving, is simply the ratio of the duration of a slip event and the interval between them: \(\bar{h} = \delta t/\Delta t\).

In the interval \(\Delta t\) between slip events, the force \(F_i\) acting on block \(i\) is slowly increased by the movement of the plate by an amount \(\Delta q\) \((k_1 \Delta q\) in dimensional units), and by the sudden movement of both neighbors by an amount \(2k_2 \Delta q\). Therefore, the force directly after the slip event is \(F_{\text{min}} = -(1+2k_2)\Delta q\) (since \(F_s = 1\)),. We can identify three ranges of the forces acting on a block:

\[
\begin{align*}
1 - (1 + 2k_2)\Delta q & \leq F_i \leq 1 - 2k_2\Delta q \quad \text{low} \\
1 - (1 + k_2)\Delta q & \leq F_i \leq 1 - k_2\Delta q \quad \text{medium} \\
1 - \Delta q & \leq F_i \leq 1 \quad \text{high}
\end{align*}
\]

A block is in the low force range after it has slipped, moves to the medium range when one neighbor has slipped, and to the high range when both neighbors have slipped. Movement within each range is caused by the slow motion of the plate.

Figure 3(a) shows a snapshot of the forces on part of the chain in the solitary regime. Peaks of only one block are present in the lower and higher force range, separated by slanted lines in the medium force range where most of the blocks reside. In Fig. 3(b) we show the distribution of forces \(P(\sigma)\) around the time of the snapshot of Fig. 3(a). \(P(\sigma)\) is peaked at \(\sigma = 1\), and \(\sigma = 1 - (1 + 2k_2)\Delta q\) due to the predominance of lines with a small slope.

The distribution of forces \(P(\sigma)\) in the solitary state shown in Fig. 3(b) is highly symmetric, hence its mean \(\bar{\sigma}\) can be approximated by the center of the distribution:

\[
\bar{\sigma} = 1 - \frac{1}{2}[1 - (1 + 2k_2)\Delta q] = \frac{1}{2}[1 - \exp(-\gamma t/\omega_0)],
\]

where we have made use of Eq. (11) for \(\Delta q\). The friction force measured in experiments is the lateral force acting on the support

\[
f = \sum\limits_{i=1}^{N} (q_i - x) = -N\bar{\sigma},
\]

where we have assumed in Eq. (7) that \(\sum(q_{i-1} - q_i) = 0\) and \(\sum q_i \approx N v_s = 0\). Since the kinetic friction force Eq. (14) is normalized by the static friction force \(F_s\), this result implies that the ratio of the kinetic to the static friction force in the solitary state can be used to extract, from experiments, the ratio \(\gamma/\omega_0\) characterizing the sliding system.

The analysis of the behavior of the forces \(\sigma_i\) in the solitary state, leads us to define a typical length scale in the system. We find that solitary motion requires two consecutive blocks in the high force range to be separated by an arbitrary number of blocks in the medium force range, and by exactly one block in the low force range. The blocks in the medium force range are arranged in monotonically increasing or decreasing slanted lines that will reach the high energy region one after the other. However, since in the time \(\delta t\) it takes block \(i\) to slip, the upper surface travels a distance \(v_s \delta t\), the absolute slope of the lines is constrained by

\[
\left| \frac{d\sigma}{dt} \right| \approx v_s \delta t.
\]

Since, in the strictly solitary state, each slanted line must start and end in the medium force zone \(\Delta \sigma = \Delta q\) wide, the minimum slope also limits the number of blocks along the line to

\[
N_{\text{line}} = \left| \frac{d\sigma}{d\sigma} \right| \Delta \sigma = \frac{\Delta q}{\Delta \sigma} = \frac{\Delta t}{\delta t}.
\]

The finite duration \(\delta t\) of a slip event introduces a typical length scale, contrary to systems displaying SOC. Strictly
speaking, in a continuous model, the size of an avalanche is given by the number of blocks performing simultaneous motion. By this definition, in a system in solitary motion, all avalanches are of size one. However, sequences up to \(N\) of size one avalanches can and do occur.

Next, we show in Fig. 4 the time evolution of \(<H_1>\) for three values of \(k_2\). We find that a transition occurs at a critical relative value of the spring constant \(k_2 \sim 1.5\). Below \(k_2\), \(<H_1>\) smoothly decreases to zero, signaling the occurrence of the solitary state, whereas, above \(k_2\), \(<H_1>\) reaches a constant finite value. An estimate of parameters for realistic sliding lubricated surfaces [10] gives \(k_2 \sim 1\). For values of \(k_2\) just below \(k_2\) there is an initial, relatively smooth decrease of \(<H_1>\), but the solitary state is reached only after many attempts. This process is shown in the left panel of Fig. 5 where we show a gray scale map of the order parameter \(<H_1>\) for \(k_2\) below and above \(k_2\). The initial uniform band corresponds to collective stick slip motion (see Fig. 2). This behavior is followed by a very short period of almost uniform motion with velocity \(v_0\), appearing as black regions in the figure. Notice that uniform motion has been shown to be unstable [2] for models with a velocity weakening friction force. Following Persson [10] we can define an energy barrier for block \(i\) as

\[
\Delta E_i = U(F_i, q_{i1}, x) - U_i(F_i, q_{i1}, x),
\]

where \(U_i(F_i, q_{i1}, x)\) is the potential energy of block \(i\), and \(U(F_i, q_{i1}, x)\) is the potential energy of the same block, moved to where it would experience the static friction force \(F_i=1\), while keeping the position of the neighboring blocks and of the plate fixed. The potential energy \(U_i\) is given by

\[
U_i = \frac{1}{2}(x - q_i)^2 + \frac{k_2}{2}(q_{i+1} - q_i)^2 + \frac{k_2}{2}(q_{i-1} - q_i)^2 + g(q_{i1}, x). \tag{18}
\]

Since \(g(q_{i1}, x)\) does not depend on \(F_i\), Eqs. (17) and (18) give

\[
\Delta E_i = \Delta E_{\text{max}}(1 - F_i^2), \quad \text{with} \quad \Delta E_{\text{max}} = 1/2\alpha_0 g. \tag{19}
\]

The probability that block \(i\) slips (i.e., overcomes the energy barrier) within a time \(dt\) is assumed to be

\[
P_i(dt) = v \exp(-\Delta E_i/k_B T)dt, \tag{20}
\]

where \(v\) is an attempt frequency, \(T\) the temperature, and \(k_B\) the Boltzmann constant. In practice, finite temperature is simulated by drawing a random number \(r_i = [0,1]\) at each integration step for each block, and if \(r_i < P_i(dt)\), where \(dt\) is the integration time step size, the static friction force is decreased to zero by setting \(h_i = 1\). In Fig. 6 the time dependence of the order parameter for collective motion \(<H_1>\) is shown at different temperatures for \(k_2 = 1\) where at \(T = 0\) the solitary state is stable. For low temperatures the order parameter goes to zero in a way similar to the zero temperature case, although small fluctuations do occur. These fluctuations grow with increasing temperature, un-
null
[13] Distance a block slips in the initial collective stick slip regime, is the maximum distance slipped by any block in the system. This distance is given by Eq. (11) with \( k_2 = 0 \), so that \( \omega = \sqrt{1 - \gamma^2} \).