Accelerated creep as a precursor of friction instability and earthquake prediction

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Earthquakes can be considered as a result of tribological instability in a system of faults of the Earth crust. Similar instabilities can be reproduced and studied in detail in laboratory-scale experiments. In this work, the earthquake model under study is a tribosystem with pronounced stick-slip behavior. Measurement of the motion of the system with a resolution of 8 nm shows that slow creep accelerated as the instability point is approached is actually observed throughout the stick stage. This motion is regular enough to serve as a basis for prediction of the onset of instability. It is shown that the motion of a solid both at the stage of slow creep and at that of fast slip is well described by the Dieterich friction law, which takes into account the dependence of friction on rate and internal state variable, if we supplement it with the contribution of local contact rigidity. In the immediate vicinity of the instability point a universal behavior is observed making possible highly accurate prediction of the onset of unstable slip from creep observations.

Keywords: static friction, sliding friction, creep, stick-slip, state variable, earthquake prediction

1. Introduction

The huge economic and social damages brought by earthquakes make their prediction an important seismological problem. The Earth crust consists of tectonic plates which move slowly relative to each other due to convective streams in the upper mantle [1, 2]. On time scales of millions of years, these motions define the structure of the Earth surface. On small time scales, these motions are responsible for earthquakes. Earthquake models are based on the fundamental observation that earthquakes for the most owe to slip along existing faults rather than to newly formed and propagating fractures in the Earth crust, and thus they are rather the subject of friction physics than that of fracture mechanics.

Since the publication of [3], it has been commonly accepted that earthquakes are stick-slip processes. A characteristic peculiarity of these processes is that fast motion accompanied by relaxation of accumulated stress begins only as the stress reaches a certain critical level. Before this level is reached, the system is in equilibrium showing no evidence of its proximity to the critical state. This peculiarity is found as well in more complex models of friction instabilities, e.g., in models of distributed tribological systems in which each element starts moving only at a certain threshold [4–8]. These systems can display a complex dynamics with known statistical properties of real earthquakes (the Gutenberg–Richer law [9–11] and Omori laws [11, 12]). The correlations in the dynamics are, however, purely statistical and can be used only for a posteriori analysis and not for prediction of a particular instability in a particular place and at a particular point in time. The response threshold of a system is a primary physical cause for the challenge of earthquake prediction. In view of the above peculiarity, some authors think that earthquakes are impossible to predict at all [13, 14].

The concept that instabilities start developing strictly after a certain threshold is, however, too simplified. Both the theory of friction and the theory of plastic deformation state that before the macroscopic stability limit is reached, creep is always observed due to microstructural stress concentrators and thermally activated deformation. From theoretical considerations it is apparent that creep rapidly speeds up in close proximity to the threshold. On this basis many noted scientists involved with mechanisms and dynamics of earthquakes consider earthquake prediction possible [1,
The question is only whether the processes of slow creep are measurable well enough and are universal to provide the basis for reliable prediction of the onset of instability.

For this question to be given an experimental answer, we studied a simple tribological model with pronounced stick-slip behavior (this model was described and studied earlier for the possibility to influence the statistics of earthquakes [20]). Leaving aside the possibility to extend the research results to real earthquakes occurring in much more complex systems, we rigidly restrict ourselves to the principle possibility to rather accurately predict the onset of instability in the simple model system. If the prediction is possible, the conducted research would form the basis for further generalization and extension of the findings to more complex distributed systems. If the prediction is impossible even in the simplest laboratory system under strictly controllable conditions, this would strengthen the position of researchers negating the earthquake predictability.

2. Experiment

The tribological system to be studied consisted of a specimen pulled along a support plate with a soft spring (Fig. 1). The specimen and spring materials were varied, but in this work we report the results only for a steel–steel pair. The coordinate of the specimen was measured using a laser vibrometer with a resolution of 8 nm. The motion of the specimen had a pronounced stick-slip character whose example is given in Fig. 2. However with a higher resolution, slow motion is actually observed throughout the stick stage, and its velocity increases as the instability point is approached (Fig. 3). Note that Figs. 2 and 3 show the same time interval, but the coordinate scales differ 1 000 times.

Figure 4 shows the time dependence of the coordinate for a series of alternate periods of “rest” and unstable slip, and Figure 5 depicts the corresponding time dependences of the velocity1. Macroscopically, the motion appears as a series of periods of full rest and slip. Actually, as noted

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1 The time dependence of the velocity from formal calculation of $\Delta v/\Delta t$ has a strong noise component due to discreteness of measurement of the coordinate. The dependences in Fig. 5 are obtained by suppression of high-frequency fluctuations of the velocity through averaging over variable intervals (depending on the proximity to the instability point). The necessity of experimental data smoothing is an important aspect of the prediction of instabilities; this aspect is considered in a separate paragraph.
above, the specimen moves throughout the stick stage; this is readily seen from the logarithmic inset in Fig. 6.

The question we tried to answer in the work is whether it is possible to describe the whole dynamics of the system both at the slow creep and fast slip stages by the same friction law. If the answer to the question was positive we could then approximate the slow creep by the universal friction law with a small number of adjustable parameters and, using this approximation, calculate the onset of unstable slip. The friction law to be studied in detail was the Dieterich friction law [21] which is primarily proposed to describe the dynamics of earthquakes.

3. Dieterich friction law

Charles Augustin de Coulomb already knew that the coefficient of static friction depends on contact time, and the sliding friction force depends on velocity. Ruina [22] in summarizing the experimental studies by Dieterich [21] in the concept of friction dependent on velocity and state variable showed that these effects are in intimate intrinsic relation. In the Dieterich friction law the friction coefficient $\mu$ depends on the instantaneous velocity $v$ and on the state variable $\theta$ as

$$\mu = \mu_0 - a \ln \left( \frac{v^*}{|v|} + 1 \right) + b \ln \left( \frac{v^\theta}{D_c} + 1 \right),$$

(1)

and the state variable is described by the kinetic equation:

$$\dot{\theta} = 1 - \frac{|v| \theta}{D_c}.$$  

(2)

The constants $a$ and $b$ in equation (1) are normally both positive and are in the range from $10^{-2}$ to $10^{-3}$; $D_c$ is the characteristic length whose value in laboratory conditions is about 1–100 $\mu$m; $v^*$ is typically ~0.2 m/s. Weak logarithmic velocity dependence (1) of the frictional force is found in tribological systems both on the macroscale [21] and on the nanoscale [23]. In the physical meaning, this dependence is governed by an exponential dependence of the rate of thermally activated processes [24–27] and is thus universal. The logarithmic velocity dependence is found for rocks, polymers, glass, paper, wood and some metals [28].

3.1. Main properties of the Dieterich friction law

In this section we briefly dwell on the main properties of the Dieterich friction law. For the state variable at rest, we have $\theta = 0$. The variable $\theta$ can thus be interpreted as the middle “age” of microcontacts. For a constant velocity $v$ and initial condition $\theta(0) = \theta_0$, Equation (2) has the solution:

$$\theta(t) = \frac{D_c}{|v|} + \left( \frac{\theta_0 - D_c}{|v|} \right) \exp \left( -\frac{|v| t}{D_c} \right).$$

(3)

Evidently the state variable $\theta$ relaxes to a new equilibrium value at a distance $D_c$. The distance $D_c$ can thus be interpreted as a characteristic slip distance over which all existing microcontacts are broken and substituted by new ones. On completion of the transient process the state variable takes a value $\theta(\infty) = D_c/|v|$, and this is also consistent with its physical interpretation as a parameter characterizing the contact time: the stationary value of $\theta$, in this case, is actually equal to the average contact time of microregularities.

For the friction coefficient in steady-state slip, we have

$$\mu = \mu_0 - (a - b) \ln \left( \frac{v^*}{|v|} + 1 \right).$$

(4)

If the slip velocity varies from $v_1$ to $v_2$, the friction coefficient varies from

$$\mu = \mu_0 - (a - b) \ln \left( \frac{v^*}{|v_1|} + 1 \right)$$

(5)

to

$$\mu = \mu_0 - (a - b) \ln \left( \frac{v^*}{|v_2|} + 1 \right).$$

(6)

For $a - b < 0$, the steady-state friction coefficient decreases with an increase in slip velocity like, e.g., in Fig. 7; this figure shows the friction coefficient measured in the steel–steel pair with a linear tribometer and slip velocity atten-
dant jump due to an abrupt tenfold increase in slip velocity at a certain point in time:

![Fig. 7. Time dependence of the friction coefficient for the steel–steel pair in experiment with the slip velocity increasing from $10^{-1}$ to $10^{-3}$ m/s at a certain point in time](image-url)
\[ \Delta \mu = (a-b) \ln \left( \frac{\nu}{\nu_0} \right). \] (7)

For the data presented in Fig. 7 it follows that \((a-b) = -0.02.\)

3.2. Creep that precedes unstable slip

At the stick stage, as shown above, slow slip (creep) actually takes place. At this stage, the system moves so slowly that the process can be considered as quasi-static. Let us consider the simplest model in which a solid of mass \(m\) is set in motion by a soft spring with a stiffness coefficient \(k\) and the free end of the spring moves with a velocity \(\nu_0\), providing that the friction force obeying the Dieterich law acts between the solid and its substrate. The equation of motion for the solid has the form:

\[ k(x_0 + \nu_0 t - x) = F_n \left( \mu_0 - a \ln \frac{\nu}{|\dot{x}|} + b \ln \frac{\nu^2}{D_c} + 1 \right). \] (8)

The latter equation in combination with the kinetic equation for the state variable

\[ \dot{\theta} = 1 - \frac{|\dot{x}| \theta}{D_c} \] (9)

completely determines the motion of the system at the creep stage. In the general case, this equation is solvable only numerically; however, ranges in which analytical estimates are possible do exist.

So for the slip velocity much lower than the steady-state creep rate: \(\nu \ll D_c/\theta_0 = \nu_0\), Equation (9) has the solution \(\theta = \theta_0\), moreover, the coordinate of the body at the left-hand side of Eq. (8) can be taken constant and equal to its initial value \(x = x_0\). From Eq. (8) we obtain

\[ |\dot{x}| = \exp \left( \frac{\mu_0 + b \ln \frac{\nu}{\nu_0} + \frac{k \nu_0 \dot{x}}{|\dot{x}|} - \frac{k \nu_0}{a F_n} - 1}{a} \right). \] (10)

On small time scales, the latter is an exponentially increasing function with a characteristic rise time of about

\[ \tau_{\nu} = \frac{a F_n}{k \nu_0}. \] (11)

If Equation (10) was valid up to the onset of instability, the instability point \(t_{stick,0}\) would be determined by the relation:

\[ \mu_0 + b \ln \frac{\nu_0}{D_c} + 1 - \frac{k \nu_0}{F_n} \] (12)

Thus, in the first approximation we have

\[ t_{stick,0} = \frac{F_n \mu_0}{k \nu_0}. \] (13)

In the next approximation, we have

\[ t_{stick,0} = \frac{F_n}{k \nu_0} \left( \mu_0 + b \ln \left( \frac{F_n \mu_0}{k \nu_0 D_c} + 1 \right) \right). \] (14)

Note that this time is much longer than \(\tau_{\nu}\). Near the instability point, function (10) has the following asymptotic behavior:

\[ |\dot{x}| = \frac{a F_n \nu^*}{k \nu_0 (t_c - t)} \] (15)

The friction coefficient reaches the maximum value

\[ \mu = \mu_0 + b \ln \left( \frac{F_n \nu_0 \nu^*}{k \nu_0 D_c} + 1 \right) \] (16)

and decreases steeply to \(\mu_0\) early in the unstable slip; its jump is thus

\[ \Delta \mu = -b \ln \left( \frac{F_n \nu_0 \nu^*}{k \nu_0 D_c} + 1 \right). \] (17)

The jump distance is determined by the relation \(\Delta x = 2 F_n \frac{|\Delta \mu|}{k}\), and hence

\[ \Delta x = 2 F_n \frac{b \ln \left( \frac{F_n \nu_0 \nu^*}{k \nu_0 D_c} + 1 \right)}{k}. \] (18)

The jump distance decreases with draw rate, and this agrees with experimental data.

In close proximity to the instability point the creep rate becomes so high that it exceeds the steady-state creep rate: \(\nu \gg D_c/\theta_0 = \nu_0\). Equation (9) thus has the form:

\[ \frac{d \theta}{dt} = -\frac{\theta}{D_c}, \quad \theta = \theta_0 e^{-t/D_c}. \] (19)

Substitution in Eq. (8) gives

\[ k F_n \left( x_0 + \nu_0 t - x \right) = \mu_0 + a \ln \frac{\nu}{\nu_0} + b \ln \frac{\nu_0}{D_c} - \frac{b \nu_0}{D_c}. \] (20)

This equation can be integrated in the explicit form:

\[ A \int_0^t \frac{\nu_0}{a F_n} \exp \left[ \frac{\nu_0}{a} (\nu_0 t - x) - \frac{\nu_0}{a} D_c \right] dt = \int_0^x \exp \left[ -\frac{\nu_0}{a} D_c \right] dx, \] (21)

where the constant

\[ A = \nu_0 \exp \left[ -\frac{\theta}{a} \right] \] (22)

is nothing but the velocity \(\dot{x}_0\) at a point in time \(t = 0\) and

\[ B = \left( \frac{b}{D_c} + \frac{k}{F_n} \right). \] (23)

The solution of Eq. (20) has the form:

\[ x = \frac{a}{B} \ln \left[ 1 - \frac{\dot{x}_0 F_n}{\nu_0} \exp \left( \frac{k \nu_0 D_c}{a F_n} - 1 \right) \right]. \] (24)

A typical creep process described by this equation is presented in Fig. 8.
For answering the question, we numerically solved the following equation of motion simultaneously with Eq. (2):

\[ m\ddot{x} = k(v_0t - x) - F_n\mu_0 - a \ln \left( \frac{\nu^*}{|x|} + 1 \right) + b \ln \left( \frac{\nu^*}{D_c} + 1 \right) \].

(29)

Here, \( m \) is the specimen mass; \( k \) is the spring stiffness; \( v_0 \) is the draw rate; and \( F_n \) is the normal force; these four parameters were determined independently. Equation (29) was solved numerically for varying parameters \( a, b, D_c, \mu_0 \) and \( \nu^* \), and the deviation between the numerical and experimental rates

\[ \delta = \int_{t_0}^{t_f} (\ln(v_{num}) - \ln(v_{exp}))^2 \, dt \]

(30)

was minimized by the steepest descent method.

Optimization over the specified parameters shows that the Dieterich friction law rather adequately describes the accelerated creep near the instability point\(^1\). However, when optimized over the range of accelerated creep and fast slip, this law underestimates the rate of very slow creep (Fig. 9). We suppose that the deviation is due to the finite contact rigidity \([11, 29]\) governed by either the specimen elasticity or the characteristic spatial scale of the microscopic potential responsible for friction\(^2\). In our experiments, the spring force increases nearly-linearly with time; hence, the finite contact rigidity is bound to provide an additional constant contribution of \( \dot{\nu} \) to the measured slip velocity. For this effect to be taken into account, we added the velocity \( \dot{\nu} \) to that found by numerically solving the equations of motion and used it as additional adjustable parameter. This allowed for a very good agreement between the theoretical and experimental dependencies over the entire velocity interval of four to five orders of magnitude (see the lower diagram in Fig. 9).

It is seen from Table 1 that \( \dot{\nu} \) is roughly proportional to the draw rate, and this supports its interpretation as the velocity associated with the contact rigidity. The characteristic interval of this velocity is, however, extremely small (from \( 10^{-5} \) to \( 10^{-7} \) m/s) and is of minor importance in the range of accelerated creep, which holds the greatest interest in our study. Figure 10 presents time dependences of the specimen velocity for several draw rates. The minimum deviation between the theoretical and experimental estimates is found for the parameters given in Table 1. The parameter \( \mu_0 \) exerts no effect on the character of motion near the instability point and determines only the time before the very first instability. The parameter \( \nu^* \) is constant from one experiment to another. The number of adjustable parameters

\(^1\)We also used the Ruina law \([22]\) as an alternative, but found that the Dieterich law describes the experiment much better.

\(^2\)These contributions, according to \([29]\), are hardly distinctive in macroscopic terms; therefore they can be modeled by the same contact rigidity.
is thus decreased to four. The obtained values of the parameters of the Dieterich law agree with the data reported in [21, 26, 28].

Note that the optimal values of the parameters $a$, $b$ and $D_c$ are in a wide spread at different draw rates and even at the same draw rate in different events of unstable slip. The spread may not be a surprise for us and is not an evidence for incorrectness of the theory. The fluctuations are physically valid because microcontacts are likely to have the only configuration throughout the creep stage. This is evident from the same order of magnitude of the total displacement as that of $D_c$ at the creep stage before the onset of instability. Thus, the creep stage lacks self-averaging of the parameters. The characteristic amplitude of their fluctuations is quite consistent with that of fluctuations of the friction coefficient in steady-state slip like, e.g., in Fig. 7. On the contrary, the parameter $\dot{\nu}$ as a representative of fast slip is stable and free of fluctuations because it is determined by averaging over many microscopic configurations.

The creep observed before unstable slip can be used to predict the onset of instability. Let us illustrate this by the example of accelerated creep immediately before the slip. In this range, the equation of motion (8) and equation (9) have analytical solution (26). It is readily seen that in this range, the ratio of the rate of the velocity $\nu$ to the acceleration $\dot{\nu}$ is determined by the relation:

$$\frac{\nu}{\dot{\nu}} = t_c - t$$  \hspace{1cm} (31)

and gives the time left before instability. Equation (31) is a linear dependence with a slope of $-1$ vanishing at the instant of a jump. Note that the linear dependence of $\nu/\dot{\nu}$ takes place with any other power dependence of the velocity on the time interval $(t_c - t)$. So, with $\nu = \text{const} \cdot (t_c - t)^{-1}$, we have $\nu/\dot{\nu} = \gamma(t_c - t)$.

Actually, the calculation of $\nu/\dot{\nu}$ for experimental data is not a trivial problem due to the presence of noise from both discreetness of the measuring system and physical inhomogeneity of slip. For the time dependences of this ratio (which are the only characteristics making possible prediction of the instability point) to be smooth, experimental data are first to be smoothed. In the next section, we discuss the problems involved in the process.

### 5. Experimental data smoothening

While the time dependence of the coordinate at the creep stage preceding the instability point is regular and has no specific peculiarity, the time dependence of the velocity does have a noticeable noise component. In accelerated creep the noise component is many times greater than its average value. Using the equations for noise-free macroscopic motion (e.g., Equation (31)) requires filtration of rapidly oscillating velocity and acceleration components unrelated to a regular process of accelerated creep. For this purpose, we are to smooth experimental data over a certain time interval which is likely bound to decrease as the instability point is approached. This brings up the problem: if the smoothing interval is too small the noise component makes the dependences of the velocity and particularly those of the acceleration absolutely irregular and their use for prediction of the instability point is thus impossible; if the interval is too large the averaging greatly “distorts” the dependences, and this also leads to false estimates of the velocity and acceleration variation. An appropriate time interval can be chosen from the following simple consideration: the averaging interval for dependences of type (10) or (25) should be shorter than the characteristic time of the velocity and acceleration variation. If the smoothed time dependence of the velocity was known, the characteristic time of the variation could be estimated as the ratio $\nu/\dot{\nu}$. Reliable prediction of the onset of instability is possible only if we manage to find an averaging interval such that the dependence of $\nu/\dot{\nu}$ is smooth enough and the averaging interval is always

<table>
<thead>
<tr>
<th>Experiment No.</th>
<th>$v_{0x}$, mm/s</th>
<th>$a$</th>
<th>$b$</th>
<th>$D_c$, µm</th>
<th>$\dot{\nu}$, m/s</th>
<th>$\mu_0$</th>
<th>$\nu^*$, m/s</th>
</tr>
</thead>
<tbody>
<tr>
<td>1</td>
<td>7</td>
<td>$4.01 \times 10^{-3}$</td>
<td>$4.75 \times 10^{-3}$</td>
<td>0.90</td>
<td>$8.48 \times 10^{-6}$</td>
<td>0.27</td>
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<td>2</td>
<td>2</td>
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<td>$8.64 \times 10^{-3}$</td>
<td>0.78</td>
<td>$9.36 \times 10^{-7}$</td>
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<td></td>
</tr>
<tr>
<td>3</td>
<td>1</td>
<td>$3.13 \times 10^{-3}$</td>
<td>$6.88 \times 10^{-3}$</td>
<td>0.70</td>
<td>$5.42 \times 10^{-7}$</td>
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</tr>
<tr>
<td>4</td>
<td>0.5</td>
<td>$3.03 \times 10^{-3}$</td>
<td>$6.20 \times 10^{-3}$</td>
<td>0.58</td>
<td>$3.21 \times 10^{-7}$</td>
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</tbody>
</table>
shorter than \( \dot{v}/\dot{\dot{v}} \) at the same point in time. For example, Figure 10 shows a typical experimental time dependence of \( \dot{v}/\dot{\dot{v}} \). In this experiment, the self-consistent condition for choosing the averaging interval is that the latter would be several times shorter than the characteristic time equal to \( \approx 0.3 \) s. This time starts considerably and systematically decreasing only in the immediate vicinity of the jump (to the onset of instability in Fig. 11 corresponds the minimum of the curve \( \dot{v}/\dot{\dot{v}} \) reached at about \( t = 3.7 \) s).

6. Universal creep in close proximity to the onset of instability

Figure 12 shows the ratio \( \dot{v}/\dot{\dot{v}} \) for the time interval beginning about 1.2 s before the instability point \( t_c \). It is seen that at the final stage, the ratio is rather adequately approximated by a linear dependence with a slope of \( -1 \), which is what we have from theoretical dependence (31). Note that
while the dependence of $\nu/\dot{\nu}$ far from the instability point is not universal, its final portion is always universal. An illustration of the universal behavior is the dependences of $\nu/\dot{\nu}$ as a function of time left before the onset of instability for several events of unstable slip at the same draw rate (Fig. 13). The dependences differ greatly far from the instability point, but display universal behavior in close proximity to it. A similar situation takes place at different draw rates (Fig. 14).

Figure 15 shows an example of the dependence of the mean-root-square error of the instability point determined from $\nu/\dot{\nu}$ as a function of time left before the instability. As the instability point is approached, the relative error decreases to about 10% and remains at this level up to the onset of instability.

7. Discussion of the results

In the context of our laboratory model we consider “long-term prediction” if the time is comparable with stick times and “short-term prediction” if it is much shorter than stick times. Our measurements show that for the tribological system under study, highly accurate short-term prediction of the onset of instability is possible if based on the universal creep preceding the instability point. For tribological systems obeying the Dietrich law, the characteristic time interval on which this universal behavior is observed is $a/\mu_0$, i.e., about 1% of the stick time. In other words, short-term prediction for the system under study is possible without any adjustable parameters in the latest stick time equal to one percent of the time of the process. At the same time, regular creep is observed during at least half the stick stage. This fact opens up a fundamental possibility of long-term prediction which, however, requires an appropriate choice of optimum parameters of the observed creep. The prediction at this stage, due to behavioral nonuniversality, will be much less accurate giving only the order of magnitude of the time left before the instability point.

Let us now discuss the possibility to extend the foregoing conclusions to earthquake prediction. An important pecu-

![Image](image1.png)

Fig. 12. Last portion of the typical dependence of $\nu/\dot{\nu}$ on the time left before instability. The portion immediately before the instability is a linear dependence that vanishes at the stability point.

![Image](image2.png)

Fig. 13. Dependence of $\nu/\dot{\nu}$ on the time left before instability for several successive instabilities at the same draw rate $\nu_0 = 1 \text{ mm/s}$.

![Image](image3.png)

Fig. 14. Dependence of $\nu/\dot{\nu}$ on the time left before instability for several draw rates: 0.5 (1), 1 (2) and 2 mm/s (3).

![Image](image4.png)

Fig. 15. Root-mean-square error of the estimated time before instability.
several spring-connected bodies are in contact with a substrate and are drawn along it with an external sliding carriage like in the theoretical model described in [4]. Clearly, each body alone is thus under the same conditions as those in experiments with a single body. Therefore, accurate measurement of the coordinate of each of the bodies is bound to allow an estimate of their proximity to the critical state. In particular, short-term prediction of the instability point for each of the bodies based on universal creep in the latest stick time of one percent is bound to remain possible.

At the same time, it is obvious that short-term prediction alone is insufficient to forecast the aftereffects of the instability for each individual body. The fact of whether unstable slip of an individual body leads to instability of the bodies surrounding it and possibly to instability of the entire system depends on a particular stress state of each body of the system, and this state can not be estimated from a mere short-term prediction for an individual body. Thus, the prediction loses most of its value, because successful prediction is not only a prediction of any instability, but also a reliable estimate of the magnitude of an event of unstable slip as a whole. Hence our experiments make clear that successful prediction of both the onset and magnitude of instability in distributed systems requires rather accurate long-term prediction of local instabilities in addition to their short-terms prediction. This problem is the subject of our laboratory research related to an upcoming project.

8. Conclusion

Thus, it is demonstrated that in the examined tribological systems, the overall specimen motion both in slow creep and in unstable slip can be adequately described by the Dieterich friction law supplemented with contact rigidity. Observations of the creep preceding the instability makes possible short-term prediction of the onset of instability and, though less accurate, its long-terms prediction. For these conclusions to be extended to real tectonic systems, more complex distributed laboratory models above all else are to be studied.

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