Meso-Scale Shear Physics in Earthquake and Landslide Mechanics

Edited by
Yossef H. Hatzor
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Cover photo: A filled bedding plane in Masada rock slopes/Photo Yossef Hatzor

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MESO-SCALE SHEAR PHYSICS IN EARTHQUAKE AND LANDSLIDE MECHANICS

Preface

This book brings together a collection of invited articles on meso-scale mechanics, as a means to develop an understanding of the underlying physics dominating shear at material interfaces and within particulate systems during conditions of rapid shearing. This edited volume emanated from the "Batsheva Seminar on Shear Physics at the Meso-scale in Earthquake and Landslide Mechanics", sponsored by the Batsheva de Rothschild fund of the Israel Academy of Sciences and Humanities and co-sponsored jointly by the US Air Force Research Laboratory and the Ben-Gurion University of the Negev.

Identification of meso-scale phenomena occurring between microscopic and continuum length scales has been one of the most exciting developments in the last decades in understanding shear between material interfaces and in particulate systems, and is considered as the bridge between the two length scales for studying material response. At the meso-scale complexities arise due to the presence of structural elements like surface roughness, grains, internal boundaries between them, physical phenomena occurring at surfaces, formation of sub-grain elements during loading, the presence of fluids, interaction between different materials, and their combined effect on the response of the system. This research area has broad applications in Geosciences and Geoengineering. For example, the initiation of seismic slip along fault planes at great depths at rates nearing shock conditions and the initiation of deep seated landslides near the earth's surface. Additionally, the basic physics of thermo-poro-mechanical coupling can be elucidated through a meso-scale mechanics approach as a means of understanding the loss of shearing resistance when water and heat are trapped inside almost impervious shear layers under great pressures. In the case of seismic slip, tremendous amount of material slips at very high velocities (several meters per second) on ultra localized shear zones. Shear heating and fluid pressurization can be associated to phenomena such as phase transition and mineral decomposition and thus play a key role in the understanding of the energetics of earthquakes. In the case of deep seated landslides, thermo-poro-mechanical processes within localized shear zones control their triggering, their sliding velocity and consequently their runout.

The 21, peer-reviewed articles are grouped into five chapters that address theoretical, computational and experimental aspects of meso-scale mechanics of material interfaces and particulate systems as follows: 1) Dynamics of frictional slip, 2) Fault gauge mechanics, 3) Experimental fault zone mechanics, 4) Granular shear and liquefaction, and 5) Dynamics of landslides.

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Yossef H. Hatzor
Jean Sulem
Ioannis Vardoulakis
Editors

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I. *Dynamics of frictional slip*

Thermo- and hydro-mechanical processes along faults during rapid slip

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ABSTRACT: Field observations of maturely slipped faults show a generally broad zone of damage by cracking and granulation. Nevertheless, large shear deformation, and therefore heat generation, in individual earthquakes takes place with extreme localization to a zone <1–5 mm wide within a finely granulated fault core. Relevant fault weakening processes during large crustal events are therefore likely to be thermal. Further, given the porosity of the damage zones, it seems reasonable to assume groundwater presence. It is suggested that the two primary dynamic weakening mechanisms during seismic slip, both of which are expected to be active in at least the early phases of nearly all crustal events, are then as follows: (1) Flash heating at highly stressed frictional micro-contacts, and (2) Thermal pressurization of fault-zone pore fluid. Both have characteristics which promote extreme localization of shear. Macroscopic fault melting will occur only in cases for which those processes, or others which may sometimes become active at large enough slip (e.g., thermal decomposition, silica gelation), have not sufficiently reduced heat generation and thus limited temperature rise. Spontaneous dynamic rupture modeling, using procedures that embody mechanisms (1) and (2), shows how faults can be statically strong yet dynamically weak, and operate under low overall driving stress, in a manner that generates negligible heat and meets major seismic constraints on slip, stress drop, and self-healing rupture mode.

1 INTRODUCTION

There has been a surge of activity in recent years towards increased physical realism in description of the earthquake process. That includes insightful geological characterization of the fine structure of fault zones, new laboratory experiments that reveal response properties in rapid or large slip, and new theoretical concepts for modeling dynamic rupture. The purpose here is to review some of those new perspectives and their impact on how we think about earthquake rupture dynamics.

1.1 *Fault zone structure, friction and a quandary in seismology*

Field observations of maturely slipped faults show a generally broad zone of damage by cracking and granulation (Chester et al., 1993), but nevertheless suggest that shear in individual earthquakes takes place with extreme localization to a long-persistent slip zone, <1–5 mm wide, within or directly bordering a finely granulated, ultracataclastic fault core (Chester and Chester, 1998; Chester et al., 2003, 2004; Heermance et al., 2003; Wibberley and Shimamoto, 2003).

On the other hand, the shear strength along a fault may be represented by

$$\tau = f\overline{\sigma} \quad \text{where} \quad \overline{\sigma} = \sigma_n - p_f \tag{1}$$

Here f is the friction coefficient, $\overline{\sigma}$ is the effective normal stress, σ_n is the total normal stress clamping the fault shut, and p_f is the pore pressure along it. It is well known that lab estimates of f

for rocks (under sliding rates of, say, μm/s to mm/s) are usually high, $f \sim 0.60$–0.85 (e.g., Byerlee (1978)).

Given that fault slip zones seem to be so extremely thin, one must conclude that if those f prevail during seismic slip, with p_f that is much closer to hydrostatic than lithostatic, we should find the following: (a) measurable heat outflow near major faults, and (b) evidence of extensive melting along exhumed faults. However, neither effect (a) or (b) is generally found.

1.2 *Weak faults, vs. statically strong faults that dynamically weaken*

There are two general lines of explanation that have been explored to resolve this quandary. One line of explanation simply postulates that major faults are *weak*. That could be because fault core materials are simply different from most rocks and have very low f, e.g., like documented for some clays and talc. Alternatively, it could be because f is not necessarily low, but pore pressure p_f is high and nearly lithostatic over much of the fault, especially down-dip where σ_n is large.

It is not the purpose here to argue against such weak-fault lines of explanation, but rather to explore an alternative which we are led to by recent observations. That is that major faults are *statically strong* but *dynamically weaken* during seismic slip. Owing to the extreme thinness of slip zones, the relevant fault weakening processes during large crustal events are likely to be thermal and, given the damage zones and geologic evidence of water-rock interactions within them, it seems reasonable to assume fluid presence. Of the various dynamic weakening processes thus far identified, it has been argued that two should be singled out as being essentially universal, in that they are expected to be active and important from the start of seismic slip in crustal events (Rice, 2006; Rice and Cocco, 2007). These are as follows:

1. Flash heating and hence shear weakening of frictional micro-asperity contacts, a process which reduces f in rapid slip, and
2. Thermal pressurization of pore fluid, which reduces the effective stress; p_f increases, because the highly granulated fault gouge is of low permeability and the thermal expansion coefficient of water is much greater than that of the rock particles.

Other thermal weakening processes may set in at large enough slip or large enough rise in temperature T_f along the fault. These include the following:

3. Thermal decomposition at large rise in T_f in lithologies such as carbonates, thus liberating a fluid product phase at high pore pressure,
4. Formation of a gel-like layer at large slip in wet silica-rich fault zones, or some related process relying on the presence of silica and water, and
5. The ultimate thermal weakening mechanism, formation of a macroscopic melt layer along the fault at large enough slip and rise of T_f, *if* the above set has not limited the actual increase of T_f to levels lower that that for such bulk melting.

While we focus on processes (1) and (2), it is very important to understand (3), (4), and (5) and others not yet identified. Still, in a sense the latter are secondary, because one expects that some significant earthquake slip, and fault weakening, will already have occurred before they can become activated. Preliminary estimates (Rempel and Rice, 2006) of when (5) would set-in suggest that with hydrostatic p_f and representative material parameter ranges, (1) and (2) are sufficiently effective at shallow fault depths that slip in significant earthquakes could often be accommodated without an onset of macroscopic melting, but that deeper in a fault zone, where the initial $\sigma_n - p_f$ (which scales the rate of heat input), and the initial T_f, are higher, melt onset should occur during increasing slip of typical surface-breaking earthquakes. For quantification of the slips and parameter ranges involved, see Rempel and Rice (2006).

Flash heating, weakening process (1), is a mechanism that has been advanced to explain high speed frictional weakening in metals (Bowden and Thomas, 1954; Archard, 1958/59; Ettles, 1986; Lim and Ashby, 1987; Lim et al., 1989; Molinari et al., 1999). It is only relatively recently that it has been considered as a process active during earthquake slip (Rice, 1999, 2006; Beeler and Tullis, 2003; Beeler et al., 2008; Tullis and Goldsby, 2003a,b; Hirose and Shimamoto, 2005;

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Noda et al., 2006; Noda, 2008). Because of the relatively low thermal conductivity of most rocks, and the relatively high shear stresses which they support at frictional micro-contacts, they are in fact susceptible to weakening by flash heating starting at sliding rates as low as 0.1 to 0.3 m/s, which is well less than the average slip rate of ~1 m/s (Heaton, 1991) inferred from seismic inversions for large earthquakes. Thermal pressurization, process (2), has independent roots in the literature on large landslides (Habib, 1967, 1975; Anderson, 1980; Voigt and Faust, 1982; Vardoulakis, 2002; Veveakis et al., 2007; Goren and Aharonov, 2009) and that on earthquakes (Sibson, 1973; Lachenbruch, 1980; Mase and Smith, 1985, 1987; Lee and Delaney, 1987; Andrews, 2002; Wibberley, 2002; Noda and Shimamoto, 2005; Sulem et al., 2005; Rice, 2006; Rempel and Rice, 2006; Ghabezloo and Sulem, 2008; Noda et al., 2009).

Process (3), thermal decomposition with generation of a high-pressure fluid phase (O'Hara et al., 2006; Han et al., 2007; Sulem and Famin, 2009) is, of course, a type of thermal pressurization. In considering process (2), the fluid phase is presumed to pre-exist in pore spaces within the fault gouge so that the pressurization begins as soon as slip and consequent frictional heating begin, whereas in (3) the fluid phase comes into existence only once enough slip, frictional heating, and temperature rise have accumulated to initiate the decomposition. Process (4) is based on findings from experiments at large but sub-seismic (in results reported thus far) slip that, in presence of water, frictional weakening at large slip is greatest for rocks of greatest silica content (Goldsby and Tullis, 2002; Di Toro et al., 2004; Roig Silva et al., 2004). The weakening is argued to be due to formation of an initially weak silica-gel layer through reaction of water with fine silica particles from fresh comminution along the shear zone. There are many studies of process (5), macroscopic melting in fault zones, of which the long-lived signature is noncrystalline pseudotachylyte veins along the fault surface and in side-wall injections. Recent contributions include Spray (1995), Tsutsumi and Shimamoto (1997), Fialko and Khazan (2004), Hirose and Shimamoto (2005), Sirono et al. (2006), and Nielsen et al. (2008).

2 DYNAMIC RUPTURE FORMULATION

2.1 *Elastodynamic methodology*

Noda et al. (2006, 2009) and Dunham et al. (2008) have begun to integrate weakening by flash heating and thermal pressurization into elastodynamic numerical methodology for spontaneous rupture development. The problems thus far addressed are of rupture along a planar fault zone within an effectively unbounded and homogeneous solid. For those, the implementation of an elastodynamic boundary integral equation (BIE), with a spectral basis set for the slip and stress distributions (Perrin et al., 1995; Geubelle and Rice, 1995) is extremely efficient and accurate. Results have been obtained for 2D anti-plane or in-plane strain. In the formulation, with the x axis passing along the fault plane, the shear stress $\tau(x, t)$ along the fault and the slip $\delta(x, t)$ are related by

$$\tau(x,t) = \tau_0(x,t) - (\mu/2c_s)V(x,t) + \phi(x,t) \tag{2}$$

where $V(x, t) \equiv \partial\delta(x, t)/\partial t$ is slip rate, μ is the shear modulus, c_s is the shear wave speed, and the functional $\phi(x, t)$ is given as a linear space-time convolution of an elastodynamic kernel, dependent on $x - x'$ and $t - t'$, with the slip $\delta(x', t')$ for all x', t' within the wave cone with vertex at x, t. Here $\tau_0(x, t)$ is some specified loading stress on the fault; it is the stress that would have been induced by the applied loadings if the fault had been constrained against any slip. We prescribe $\tau_0(x, t)$ as a uniform background stress τ^b for all time, plus some localized overstress applied at $t = 0$ to nucleate rupture. $\delta(x, t)$ and $\phi(x, t)$ are expanded in a Fourier basis set, so the convolution is expressed by

$$\begin{Bmatrix} \delta(x,t) \\ \phi(x,t) \end{Bmatrix} = \sum_{n=-N/2}^{N/2} \begin{Bmatrix} D_n(t) \\ \Phi_n(t) \end{Bmatrix} \exp(in\hat{k}x) \quad \text{with } \Phi_n(t) = \int_0^t C_n(t-t')D_n(t')dt' \tag{3}$$

Here N is a large even integer, the $C_n(t)$ are known real functions (Perrin et al., 1995; Geubelle and Rice, 1995), $\hat{k} = 2\pi/X$ where the periodic repeat length X of the truncated Fourier series is chosen large enough that waves from the periodic replications of the rupture event do not arrive to neighboring replications in the time of interest, and the $D_{-n}(t)$ are complex conjugates of the $D_n(t)$ with D_0 and $D_{N/2}$ being real. Eq. (3) is equivalent to a real Fourier cosine and sine series truncated at N terms (with no sine term when $n = N/2$). Through FFT procedures, the histories $D_n(t)$ are determined by $\delta(x,t)$ at N sample points, equally spaced by $\Delta x = X/N$, i.e., by the histories $\delta(j\Delta x - X/2, t)$ for $j = 0, 1, 2, \ldots, N-1$, and conversely. That greatly speeds calculations. The $\Phi_n(t)$ and $\phi(j\Delta x - X/2, t)$ are similarly related.

2.2 *Friction law with weakening by flash heating and thermal pressurization*

The simple flash heating model reviewed here (Rice, 1999, 2006; Beeler and Tullis, 2003; Beeler et al., 2008) was intended to approximately determine the expression for f of eq. (1) for conditions of sustained sliding at some speed V. The f so derived must be regarded as a *steady state* value, written here as f_{ss} and is regarded as a function of slip rate, $f_{ss} = f_{ss}(V)$, although it is also a function of the spatially averaged (over patches of fault area large enough to include many contacts) temperature T_f of the fault plane, which evolves with ongoing slip and time. In the model (Rice, 2006) it is assumed that contact temperature $T_c = T_f$ when a contact pair first forms, but then as the contact slides during its brief lifetime (which is D/V for a contact asperity of diameter D; see Fig. 1), T_c rises substantially above T_f. The rise is due to the intense localized heat generation at the contact, at rate $\tau_c V$, where τ_c is the contact shear strength (typically of order 0.1 times shear modulus μ at low T_c; see discussion in Rice (2006)). τ_c is assumed to have negligible variation as T_c increases, but then to abruptly decrease to a much lower weakened value τ_w when T_c reaches a "weakening" temperature T_w. Within the model, based on 1D heat conduction at the sliding contact like in Archard (1958/59), the slip rate such that an asperity of diameter D would begin to weaken only just as it is slid out of existence is then (Rice, 1999, 2006)

$$V_w = (\pi\alpha_{th}/D)(\rho c(T_w - T_f)/\tau_c)^2 \tag{4}$$

Here ρc is volumetric specific heat and $\alpha_{th} = K/\rho c$ is thermal diffusivity (K is thermal conductivity). Estimates based on rock properties and assumed D of order 10 μm, as well as comparison to experimental results of Tullis and Goldsby (2003a, b) and Beeler et al. (2008) (by rotary shear of a rock annulus in an Instron frame), suggest that V_w is of order 0.1 to 0.3 m/s for rocks such as quartzite, feldspar, granite and gabbro when T_f = room temperature. V_w is expected to be less at higher T_f. Thus the model takes contacts to be strong for all their lifetime at low slip rates, $V < V_w$, but to be strong for only a fraction of their lifetimes at high slip rates, $V > V_w$, that fraction being V_w/V.

Neglecting the actual statistical distribution of contact diameters, and taking D as a representative value, these concepts lead to the steady-state friction coefficient

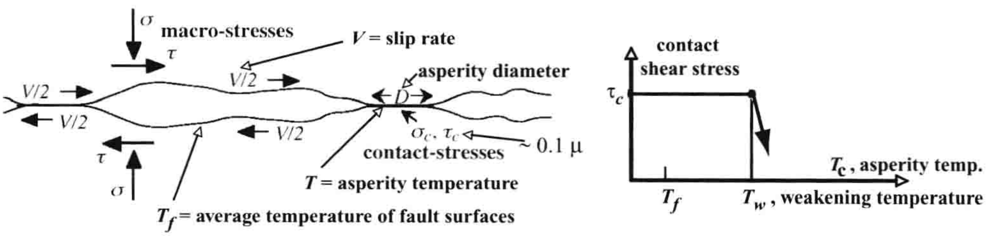

Figure 1. Simple representation of asperity contacts and their strength loss for flash heating model.

$$f_{ss}(V) = \begin{cases} f_0 & V \leq V_w \\ f_w + (f_0 - f_w)V_w/V & V \geq V_w \end{cases} \tag{5}$$

where f_0 is the low-speed friction coefficient and $f_w = f_0 \tau_w / \tau_c$ is the value to which it would reduce if all contacts were in the weakened state. Note that despite the notation $f_{ss}(V)$, the dependence mentioned on ambient fault temperature enters from the dependence of V_w on T_f, eq. (4). Rough estimates from the Tullis and Goldsby (2003a, b) experiments, which covered a limited velocity range, up to slightly less than $V = 0.4$ m/s, are that $f_0 \approx 0.64$ and $f_w \approx 0.12$ for quartzite, $f_0 \approx 0.82$ and $f_w \approx 0.13$ for granite, and $f_0 \approx 0.88$ and $f_w \approx 0.15$ for gabbro. However, the f_w involve significant extrapolation and, in experiments of Yuan and Prakash (2008) on quartzite from the same source and annular configuration, but in a Kolsky bar dynamic torsion apparatus, f was only very slightly below 0.20 at slip rates as high as 2–4 m/s. Eq. (5) and the small $V_w \sim 0.1$ m/s for quartzite, suggest that f should nearly coincide with f_w at such rates, thus that $f_w \approx 0.18$–0.20.

In fact, we cannot simply assume f to be a decreasing function of V in eq. (1) because that makes the problem of sliding between elastic continua ill-posed (there is a short wavelength divergence in response to small initial perturbations from steady sliding). The problem is remedied (Rice et al., 2001) mathematically when we look to experiments and embed the description of variations of f in rate and state friction concepts. Thus, with a "slip" version of state evolution, f is assumed to be given by the form (Rice, 1983)

$$\frac{df}{dt} = \frac{a}{V}\frac{dV}{dt} - \frac{V}{L}(f - f_{ss}(V)) \tag{6}$$

where a is the direct effect coefficient in rate and state modeling ($a \approx 0.01$ at room T for quartzite and granite, and it scales in approximate proportion to absolute T), and L is a slip distance adequate to renew the asperity contact population, typically taken as 5–20 μm in our studies, as guided by observed state evolution slip distances in rate and state friction experiments. The smallness of L makes the simulations extremely challenging in terms of present-day computers and, as of the writing, the longest rupture lengths simulated are ~30 m (Dunham et al., 2008; Noda et al., 2009).

In the numerical simulations (Noda et al., 2009; Dunham et al., 2008) we replace the constant term f_0 in eq. (5) with $f_{LV}(V)$, which is the weak logarithmic function of V describing slow-rate

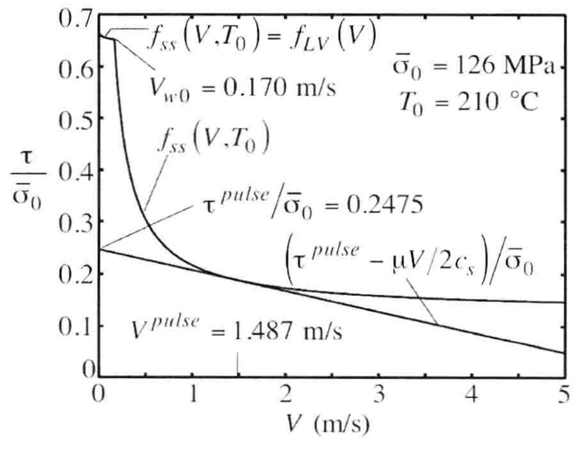

Figure 2. From Noda, Dunham, and Rice (2009), their figure 1. Plot of f_{ss} based on Tullis and Goldsby (2003a, b) parameters for granite, adjusted to conditions at mid-seismogenic depth, ~7 km, for crustal strike slip earthquakes. Plot is based on assumed ambient effective normal stress $\bar{\sigma}_0 = 126$ MPa and temperature $T_0 = 210$°C; as slip develops, heating increases T_f over T_0, and pore fluid pressurization will reduce $\bar{\sigma}$ from $\bar{\sigma}_0$. The Zheng and Rice (1998) stress level τ^{pulse}, important to understanding whether rupture takes the self-healing pulse or classical shear crack mode, is also shown.

friction at steady-state. That is, $f_{LV}(V) = f_0 - (b - a)\ln(V/V_0)$ in the standard notation, where $b - a \approx 0.002$–0.004 for rate-weakening frictional surfaces in granite. Here V_0 can be chosen arbitrarily in that low-speed regime (say, as 1 μm/s) with corresponding adjustment of f_0 to correspond to the correct f_{ss} at that rate. Fig. 2, taken from Noda et al. (2009), plots the resulting f_{ss} using parameters which are thought to be representative for granite, at conditions corresponding to a mid-seismogenic zone depth for crustal strike-slip earthquakes. Also, while irrelevant in our simulations, for some purposes the a/V term in eq. (6) should be regularized near $V = 0$ in a manner consistent with the thermal activation basis for the direct effect (Rice et al., 2001).

2.3 *Inclusion of thermal pressurization of pore fluid in friction formulation*

In order to close our system of governing equations which, so far, consists of eqs. (1), (2), (3) and (6), it is necessary to relate the histories of p_f and T_f to that of slip and stressing along the fault. That is because p_f directly enters eq. (1) and T_f is a parameter on which $f_{ss}(V)$ depends (see eqs. (5) and (4)). The new ingredients which close the system are the equations of conservation of energy, with inclusion of conductive heat transfer, and of conservation of fluid mass with inclusion of Darcy fluid seepage and poro-thermo-elastic considerations. These amount to a pair of coupled PDEs for the fields of pore pressure p and temperature T near and on the fault. The values of p and T thus determined at the fault surface itself are the respective p_f and T_f that we seek.

In writing the conservation laws we neglect certain apparently minor terms (e.g., advective heat transfer by moving fluid), and recognize that the gradients of pore pressure p and temperature T very near the fault are generally very much larger in the direction perpendicular to the fault (the z direction) than in the x or y directions which are parallel to it (such may sometimes not be a valid assumption immediately at the moving rupture tip). Thus (e.g., Rice (2006)) we have

$$\rho c \frac{\partial T}{\partial t} = \frac{\partial}{\partial z}\left(\rho c \alpha_{th} \frac{\partial T}{\partial z}\right) + \tau \frac{\partial \gamma^{pl}}{\partial t} \quad \text{and} \quad \beta\left(\frac{\partial p}{\partial t} - \Lambda \frac{\partial T}{\partial t}\right) + \frac{\partial n^{pl}}{\partial t} = \frac{\partial}{\partial z}\left(\beta \alpha_{hy} \frac{\partial p}{\partial z}\right) \tag{7}$$

Here $\dot{\gamma}^{pl}$ is the inelastic fault-parallel shear strain rate and $\dot{n}^{pl}$ is the inelastic dilatancy rate (n itself is the volume of pore space per unit aggregate volume of porous material, that aggregate volume being measure in some reference state before the deformation episode considered). Also, ρc is the specific heat per unit volume, $\alpha_{th} = K/\rho c$ is the thermal diffusivity, and K is thermal conductivity; β is a porous medium storage coefficient under the particular mechanical constraints near a fault zone (see Rice (2006)), $\alpha_{hy} = k/\beta\eta_f$ is the hydraulic diffusivity, k is permeability, and η_f is viscosity of the pore fluid; and Λ is a parameter representing dp/dT due to heating under undrained, elastically reversible conditions. Rice (2006) and Rempel and Rice (2006) compile estimates of these various parameters at mid-seismogenic depths in the crust, based principally on data of Wibberley (2002) and Wibberley and Shimamoto (2003) for gouge of the Median Tectonic Line Fault (Japan) under a range of confining stresses, and on tabulated thermophysical data for water and minerals.

Rempel and Rice (2006) also compare fully non-linear solutions of eqs. (7) to the linearized versions

$$\frac{\partial T}{\partial t} = \alpha_{th} \frac{\partial^2 T}{\partial z^2} + \frac{\tau}{\rho c} \frac{\partial \gamma^{pl}}{\partial t} \quad \text{and} \quad \frac{\partial p}{\partial t} - \Lambda \frac{\partial T}{\partial t} + \frac{1}{\beta} \frac{\partial n^{pl}}{\partial t} = \alpha_{hy} \frac{\partial^2 p}{\partial z^2} \tag{8}$$

(with coefficients ρc, α_{th}, β and α_{hy} considered constant) for cases with *a priori* specified histories of $\partial \gamma^{pl}/\partial t$ and $\partial n^{pl}/\partial t$, e.g., representing earthquake slip at a specified constant rate in time, with the aim of estimating (using eq. (1) but assuming a constant low f as motivated by flash heating) the relation between stress τ and slip δ during seismic rupture. They find that the procedure adopted by Rice (2006), of iteratively choosing the constant values of coefficients in the linearized PDEs of eqs. (8), as certain path averages in p, T space of those same coefficients, when regarded as known functions of p, T, along the p, T path predicted by eqs. (8), gives a tolerable match to results for τ *versus* δ based the full non-linear solutions to eqs. (7). In our spontaneous elastodynamic analyses presented here we have used eqs. (8), without that iterative choice of coefficients, but rather with coefficients based on the ambient p, T at the mid-seismogenic depth ($\sim$7 km) considered.

The conceptually (but not computationally) simplest case for eqs. (7) or (8) is that of slip on a mathematical plane. In that case $\partial\gamma^{pl}/\partial t = V(x,t)\delta_{\mathrm{Dir}}(z)$, where V is the local slip rate and $\delta_{\mathrm{Dir}}(z)$ is the Dirac function. However, actual shear zones, when highly localized to a prominent slip surface, can nevertheless distribute deformation over regions that may extend over a few 10s to a few 100s of μm. Because the difference between such small but finite thicknesses and zero thickness is sometimes not negligible on the seismic time scale in our modeling, a Gaussian distribution of shear heating $\tau\,\partial\gamma^{pl}/\partial t$, like in Andrews (2002), over a zone of nominal thickness $2w$ is assumed. We write

$$\tau\frac{\partial\gamma^{pl}}{\partial t} = \frac{\tau(x,t)V(x,t)}{\sqrt{2\pi}w}\exp\left(-\frac{z^2}{2w^2}\right) \tag{9}$$

which reduces to $\tau\,\partial\gamma^{pl}/\partial t = \tau(x,t)V(x,t)\delta_{\mathrm{Dir}}(z)$ as $w \to 0$.

In the numerical studies (Noda et al., 2008; Dunham et al., 2008) of spontaneous rupture based on this formulation (which have thus far taken $\dot{n}^{pl} = 0$), explicit finite difference (in z and t) versions of eqs. (8) are solved at each elastodynamic gridpoint location $x_j = -X/2 + jX/N$ along the rupture, ultimately to give, with help of the other governing equations, $T_f(x_j,t)$ and $p_f(x_j,t)$. For what we think to be appropriate ranges $\alpha_{th} \sim 1$ mm^2/s and $\alpha_{hy} \sim 1$–10 mm^2/s (Rice, 2006; Rempel & Rice, 2006), it turns out that when the elastodynamic time steps Δt are already short enough to resolve the state evolution of eq. (6) with $L \sim 5$–$20\,\mu$m, explicit finite difference solution of eqs. (8) requires significantly shorter time steps, to meet the requirement is that diffusion grid spacing be sufficiently small that the error associated with the spatial discretization of the diffusion equations is comparable the error in the elastodynamic system. Thus Noda et al. (2008) devised a procedure based on a quadratic interpolant of $\phi(x,t)$ within an elastodynamic time step Δt, to achieve second-order accuracy. The interpolant is constructed from $\phi(x,t-\Delta t)$, $\phi(x,t)$, and $\phi(x,t+\Delta t)$ for use between t and $t+\Delta t$, so that the condition of eq. (2), $\tau = \tau_0 - (\mu/2c_s)V + \phi$, along with eq. (1), eq. (6) with eqs. (4) and (5) (with f_0 replaced by $f_{LV}(V)$), and eqs. (8) are satisfied, to numerical accuracy, within each of the smaller time steps for diffusion. This is highly accurate but very demanding computationally because of the small but presumably realistic $L \sim 5$–$20\,\mu$m and realistic diffusivities, $\sim$1–10 mm^2/s, that we use.

2.4 *Theoretical background on strong rate-weakening and self-healing slip pulses*

The adopted friction description involves strong rate-weakening (see Fig. 2). To provide background for understanding when strong rate-weakening will lead to rupture in the mode of a self-healing slip pulse, versus a classical enlarging shear crack, we digress here to review results from Zheng

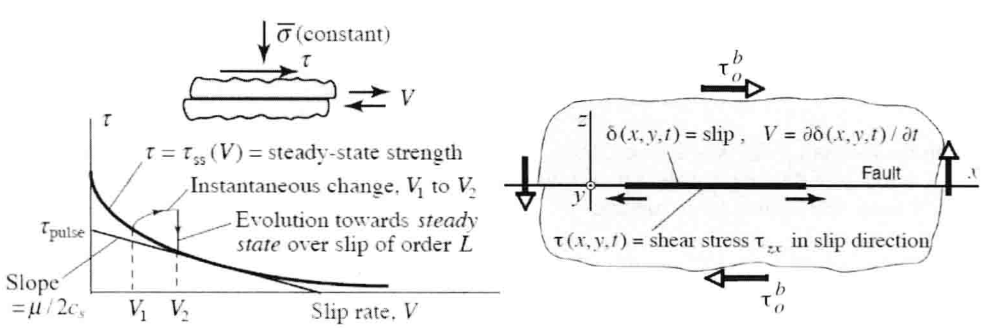

Figure 3. Friction with strong rate dependence. It is assumed here, to simplify, that the steady-state friction $f_{ss}(V)$ is a function of slip rate V only, and that the effective normal stress $\bar{\sigma}$ is constant; then the steady-state shear stress $\tau_{ss}(V) = f_{ss}(V)\bar{\sigma}$. (Modified from Rice (2001).)

and Rice (1998), in part following a recapitulation in Rice (2001), discussed with reference to Fig. 3. Some representative results from Dunham et al. (2008) and Noda et al. (2009), based on the formulation in the earlier parts of this Section 2, are presented in Section 3 to follow.

Studies by Cochard and Madariaga (1994, 1996), Perrin et al. (1995), and Beeler and Tullis (1996) had shown that strong rate-weakening could lead to rupture in the self-healing mode. The objective in Zheng and Rice (1998) was then to establish guidelines for when this type of rupture mode would in fact result. To simplify enough that sharp results could be established, it was assumed in that development that the steady state shear strength $\tau_{ss}(V)$ was a function *only* of slip rate V, which could, e.g., result if $\overline{\sigma}$ is constant and $f_{ss}(V)$ is a function only of V. Then, $\tau_{ss}(V) = f_{ss}(V)\overline{\sigma}$. However $\tau_{ss}(V)$ does not give the expression for τ for variable V, because $\tau = f\overline{\sigma}$ and f evolves according to eq. (6). Thus the response to a sudden change in V is as depicted on the left in Fig. 3, a response which also assures well-posedness in problems of frictional sliding between elastic continua (Rice et al., 2001).

Now consider a fault surface which we treat as the boundary $z = 0$ between two identical half spaces (Fig. 3, at right). An initial shear stress $\tau_0(x, y) = \tau_0^b$, a constant level too small to cause failure, acts everywhere on S_∞ (the entire x, y plane) except in small nucleation region S_{nucl} which will be overstressed to start the rupture. The stress level τ_{pulse} (the same as what Noda et al. (2009) denoted τ^{pulse} in their diagram shown here in Fig. 2) is marked in Fig. 3. τ_{pulse} is defined as the largest value that τ_0^b could have and still satisfy $\tau_0^b - \mu V/2c_s \leq \tau_{ss}(V)$ for all $V > 0$.

Suppose that $\tau_0^b < \tau_{\text{pulse}}$. As will now be seen, that effectively precludes the possibility that rupture could occur on S_∞ in the form of an indefinitely expanding shear crack. Note that

$$\tau_0^b < \tau_{\text{pulse}} \text{ implies that } \tau_{ss}(V) - (\tau_0^b - \mu V/2c_s) > 0 \quad \text{for all } V > 0. \tag{10}$$

Use is made of an elastodynamic conservation theorem (Zheng and Rice, 1998)

$$\int\int_{S_\infty} [\tau(x, y, t) - \tau_0(x, y) + \mu V(x, y, t)/2c_s]\, dxdy = 0 \tag{11}$$

which holds throughout the rupture; $\tau_0(x, y)$ is the stress field at $t = 0$ when rupture is nucleated.

Assume that, with $\tau_0^b < \tau_{\text{pulse}}$, rupture has been locally nucleated and grows on S_∞ in the form of an *indefinitely expanding shear crack*. Such an assumption can be shown, as follows, to lead to a definite *contradiction* in the case of mode III (anti-plane) slip, and to seemingly implausible result in general, meaning that we must reject the assumption that an indefinitely expanding crack-like rupture is possible when $\tau_0^b < \tau_{\text{pulse}}$. To see why, note that the integrand in eq. (11) everywhere along the rupturing surface $S_{\text{rupt}}(t)$ where $\tau \approx \tau_{ss}(V)$, except for S_{nucl} and for small regions at the rupture front affected by the rate/state regularization (so that τ may depart significantly from $\tau_{ss}(V)$ in those regions; they are small because L is in the range of a few to a few tens of microns; see discussions in Zheng and Rice (1998) and Noda et al. (2009) for further quantification), is equal to

$$\tau_{ss}(V) - \tau_0^b + \mu V/2c_s = \tau_{ss}(V) - (\tau_0^b - \mu V/2c_s) > 0 \tag{12}$$

where the inequality follows from eq. (10). Thus, denoting by $S_{\text{out}}(t)$ ($= S_\infty - S_{\text{rupt}}(t)$) the region of S_∞ lying outside the rupture at time t, and noting that $V = 0$ there, we must by eq. (11) then have, for any sufficiently large rupture,

$$\Delta F_{\text{out}}(t) \equiv \int\int_{S_{\text{out}}(t)} [\tau(x, y, t) - \tau_0^b]\, dxdy < 0 \quad (\text{and } \Delta F_{\text{out}}(t) \to -\infty \text{ as } S_{\text{rupt}}(t) \to \infty) \tag{13}$$

where $\Delta F_{\text{out}}(t)$ is the change in total shear force (positive in the direction of initial shear stressing) supported *outside* the ruptured zone. Inequality (13), however seems implausible: We expect ruptures to result in an increase $\Delta F_{\text{out}}(t)$ in the net force carried outside themselves, or at least to

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