Acoustic fluidization for earthquakes?

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Abstract: Melosh [1996] has suggested that acoustic fluidization could provide an alternative to theories that are invoked as explanations for why some faults appear to be weak. We show that there is a subtle but profound inconsistency in the theory that unfortunately invalidates the results. We propose possible remedies but must acknowledge that the relevance of acoustic fluidization remains an open question.
1 Introduction

In the standard rebound theory of earthquakes, deformation elastic energy is progressively stored in the crust and is suddenly released in an earthquake when a threshold is reached. The Ruina-Dieterich friction laws [Dieterich, 1972; 1978; Ruina, 1983] constitute the basic ingredient used to describe the interaction between the two sides of the sliding fault. Friction coefficients based on laboratory experiments [Byerlee, 1977; Scholz, 1998] fail to account for modern observations of strain [Jackson et al., 1997], stress [Zoback et al., 1987; Zoback, 1992a; 1992b] and heat flows [Heney and Wasserburg, 1971; Lachenbruch and Sass, 1980; 1988; 1992; Lachenbruch et al., 1995] (see Sornette [1999] for a synthesis).

Resolutions of these paradoxes usually call for additional mechanisms, involving fluids [National Research Council, 1990], crack opening modes of slip [Brune et al., 1993], dynamical collision effects [Lomnitz, 1991; Pisarenko and Mora, 1994], frictional properties of a granular gouge model under large slip [Scott, 1996], space filling bearings with compatible kinematic rotations [Herrmann et al., 1990], hierarchical scaling [Schmittbuhl et al., 1996], etc.

Melosh [1996] has recently suggested that the mechanism of “acoustic fluidization” could provide an alternative to theories that invoke pressurized fluids as an explanation for why some faults appear to be weak. Fluidization usually refers to the experimental observation that granular material in the presence of an interstitial fluid can liquidify when shaken sufficiently strongly [Russo et al., 1995]. The liquifaction is due to the fact that granular media are first compressive for small deformation leading to an increase of the interstitial fluid pressure. This increase in turn decreases the friction between the grains that can eventually become free to shear. For example, liquifaction of sediments by resonant amplified seismic waves have been proposed to be in part responsible for the damage and collapse of certain buildings during the Michoacan earthquake, 1985 [Lomnitz, 1987] and for the damage in the Marina district of San Francisco during the Loma Prieta earthquake [Bardet, 1990].

In the acoustic fluidization mechanism [Melosh, 1979; 1996], no interstitial fluid is invoked. A fraction $e$ of the earthquake energy is released as high-frequency acoustic waves that scatter off and shake the granular gouge leading to the build-up of a
local acoustic pressure. According to [Melosh, 1996], when this pressure becomes of the order of the overburden lithostatic pressure \( \rho gh \), the granular gouge becomes essentially free to slip without much residual friction.

The purpose of this note is to show that there is a problem with this mechanism because it predicts a slip velocity during an earthquake more than two orders of magnitudes smaller than the typical meters per second for observed earthquakes, in contradiction to the result of Melosh [1996]. The problem stems from a confusion in the definition of dissipation and scattering lengths. We then suggest possible modifications of Melosh’s theory that could resolve this problem and which lead to a richer theory.

2 Summary of Melosh’s theory and useful background

2.1 Acoustic wave energy transport

The first ingredient is the generation and transport of high-frequency acoustic waves in the core of the fault. Melosh [1996] uses the standard diffusion equation (his equation 2) for the elastic transport of acoustic waves [Ishimaru, 1978; Sornette, 1989a-c] with a dissipation and a source term:

\[
\frac{dE}{dt} = \frac{\xi}{4} \nabla^2 E - \frac{c}{\lambda Q} E + e \dot{\tau} ,
\]

where \( E \) is the acoustic wave elastic energy density. The diffusion coefficient \( \xi/4 \), where \( \xi \) is called the scattering diffusivity by Melosh [1996], can be expressed in terms of the elastic mean free path \( l_e \) and of the transport acoustic wave velocity \( c \) at scales below \( l_e \). The velocity \( c \) is of the order of the shear wave velocity [Turner and Weaver, 1996; Van Albada et al., 1991; Van Tiggelen and Lagendijk, 1993]. We thus get [Ishimaru, 1978; Sornette, 1989a,c] :

\[
\frac{\xi}{4} \approx \frac{1}{3} c l_e .
\]

We stress that the term “elastic” refers to the fact that \( l_e \) is the characteristic distance over which an acoustic wave propagates before being scattered in other directions, without any loss of energy.
The l.h.s. and first term of the r.h.s. of (1) give the diffusion equation which describes the transport of an acoustic wave in a multiple-scattering medium. The second term of the r.h.s. of (1) will be shown to describe the presence of a genuine absorption, while the last source term corresponds to the conversion to acoustic waves of a fraction $e$ of the mechanical work performed per unit time by the fault motion with strain rate $\dot{\varepsilon}$ and shear stress $\tau$.

### 2.2 Diffusive transport

The first two terms of (1)
\[
\frac{dE}{dt} = \left( \frac{\xi}{4} \right) \nabla^2 E
\]
gives the standard parabolic diffusion equation which is based on the following processes. Once generated from a source, an acoustic wave propagates roughly ballistically over a typical distance of the order of the elastic (scattering) mean free path $l_e$. Over this distance, the equation governing the acoustic wave propagation is the hyperbolic wave equation
\[
\frac{\partial^2 A}{\partial t^2} = c^2 \nabla^2 A ,
\]
where the wave amplitude $A$ is related to $E$ by $E = |A|^2$ and the wave velocity $c$ may depend locally on position to reflect the heterogeneity of the medium. Due to this heterogeneity, the wave is scattered off its initial propagation path along the direction $x$ and its intensity in this direction $x$ decays as $\exp[-x/l_e]$. This exponential decay of the intensity does not correspond to a genuine absorption but rather reflects the loss of acoustic energy along the direction $x$ to all possible scattered waves in all other directions. Mathematically, the exponential decay $\exp[-x/l_e]$ can be derived from (4) using standard scattering theory [Ishimaru, 1978]. The conservation of acoustic energy is ensured by the fact that the sum of wave intensity over all directions of propagation remains constant.

Beyond the distance $l_e$, the nature of the transport of the wave intensity crosses over from ballistic (i.e. straight propagation) to diffusive, corresponding to the picture where the acoustic wave can be viewed as a superposition of random walks with typical step length equal to $l_e$. This means that the equation for the wave propagation changes from the hyperbolic wave equation for the wave amplitude to the parabolic diffusion equation for the wave intensity given by the first two terms of (1).
One can quantify this by the following example. Consider an acoustic wave of energy $E_0$ impinging on a slab of thickness $L$ made of heterogeneities that scatter off the acoustic wave, and whose scattering strength is quantified by the elastic mean free path $l_e$. Anderson [1985] and Sornette [1989c,d] have revisited this diffusion equation to get the transmission coefficient in this example, i.e. the acoustic energy which is transmitted to the other size of the slab, as a function of its thickness $L$. The result is:

$$E(L) \simeq E_0 \frac{l_e}{L}.$$  \hspace{1cm} (5)

Note that the decay follows the algebraic $1/L$ law rather than an exponential law. Furthermore, the acoustic intensity profile within the slab is linear and not exponentially decreasing:

$$E(z) \simeq E_0 \frac{L + l_e/3 - x}{L}, \quad \text{for } l_e < x < L - l_e.$$  \hspace{1cm} (6)

These results (5,6) highlight that the diffusive transport of the acoustic energy due to multiple scattering event is very different from the exponential attenuation that a genuine absorption would produce.

### 2.3 Absorption

The third term $-(c/\lambda Q)E$ of (1) quantifies genuine absorption processes. The parameter $\lambda$ is the acoustic wavelength and $Q$ is the quality factor. To see that this term reflects absorption, we consider (1) in absence of the spatial derivative $\nabla^2 E$ and of the last source term:

$$\frac{dE}{dt} = -\frac{c}{\lambda Q} E.$$  \hspace{1cm} (7)

Its solution is

$$E(t) = E_0 \exp \left( -\frac{c}{\lambda Q} t \right),$$  \hspace{1cm} (8)

which is very different from the energy decay (5) solely due to diffusion. It is thus clear that the term $-(c/\lambda Q)E$ is not coming from elastic scattering but solely from genuine absorption, i.e. conversion of acoustic energy into thermal energy.

The usual definition of the quality factor $Q$ is [Knopoff, 1964]

$$Q \equiv 2\pi \frac{l_a}{\lambda},$$  \hspace{1cm} (9)
where \( l_a \) is the absorption length defined by the exponential decay \( \exp(-x/l_a) \) of a *ballistically* propagating wave in an absorbing medium.

Melosh [1996] introduces a characteristic length \( l_* \), which he calls (misleadingly) the “scattering length”, defined by

\[
l_* \equiv \sqrt{\frac{\xi \lambda Q}{4c}}. \tag{10}
\]

Using (2) and (9), we get

\[
l_* \approx \sqrt{\frac{2\pi}{3}} \sqrt{l_e l_a}. \tag{11}
\]

Calling \( l_* \), a “scattering length”, is misleading because \( l_* \) is in reality the effective absorption length in the diffusive medium. To see this, we use the standard diffusion relation

\[
l_*^2 \approx 6\frac{\xi}{\tau_a} \tag{12}
\]

linking the radius of gyration \( l_* \) covered by a diffusing process over a time \( \tau_a = l_a/c \) equal to the time needed for the wave to cover the real distance \( l_a \), along its convoluted multi-scattered path. The prefactor 6 holds for diffusion in a three dimensional space.

Using (2), we get \( l_* \approx \sqrt{2l_e l_a} \), which recovers the (11) up to a numerical factor of order unity. The expression (11) can be derived by several other methods [Sornette, 1989d]. What is important is that \( l_* \) scales as the geometrical mean of \( l_e \) and \( l_a \), which comes from the random walk nature of the diffusive process. Physically, in the diffusive regime, the acoustic wave energy is absorbed over the characteristic length \( l_* \), which stems from the fact that, to cross the distance \( l_* \), the wave follows random walk paths of length \( l_a \sim l_*^2/l_e \). This reflects that attenuation of a wave in a scattering medium is a function of both absorption of energy and scattering.

### 2.4 Feedback of the acoustic vibrations on the slip rate

The interesting idea of Melosh [1996] is that the high-frequency vibrations may shake the fault and unlock it, leading to an easier sliding motion. For this, he proposes the following effective friction equation, relating the strain rate \( \dot{\varepsilon} \), the shear stress \( \tau \) and the normalized acoustic wave energy \( \Psi = E \rho c^2/(\rho gh)^2 \):

\[
\dot{\varepsilon} = \frac{\tau}{\rho \lambda c} \left[ \frac{1 - \text{erf}(\frac{1}{2\sqrt{\Psi}})}{1 + \text{erf}(\frac{1}{2\sqrt{\Psi}})} \right]. \tag{13}
\]
The main physical phenomenon taken into account in this equation is that, due to the acoustic shaking, the effective viscous friction $\tau/\dot{\epsilon}$ is a decreasing function of the acoustic wave energy. This mechanism is related to the velocity weakening mechanism induced by collision between asperities that lead to a transfer of momentum from the direction parallel to the fault to the direction transverse to it [Lomnitz-Adler, 1991; Maveyraud et al., 1998].

Putting (13) in (11) and looking for stationary modes gives the non-linear ordinary differential equation (21) (for the case $\eta = 1$, see below), whose analysis leads to the prediction of two rupture modes [Melosh, 1996].

3 Problem with Melosh’s theory

In order to obtain realistic values, there are some constraints that the model parameters must satisfy. The key parameter is the “regeneration” parameter

$$r \Sigma^2 = \frac{eQ}{2} \left( \frac{\tau}{\rho gh} \right)^2,$$

(14)

where $\Sigma = \tau/\rho gh$ is the normalized shear stress. Melosh [1996] finds reasonable solutions only for $r \Sigma^2 > 2.8$. For a typical fraction $e \approx 0.1$ of conversion to acoustic waves of the mechanical work performed per unit time by the fault motion and for a ratio $\tau/\rho gh$ as low as 0.1 as suggested from observations on the San Andreas fault, this value $r \Sigma^2 > 2.8$ corresponds to $Q > 5600$. This estimation may vary by an order of magnitude with the conversion factor and the relative shear stress. However, the message is that the quality factor $Q$ measuring the attenuation of the acoustic waves must be high, in the range of $10^3$ for the acoustic waves to be self-sustained during the earthquake slip motion. This is the first condition.

On the other hand, Melosh’s theory predict the slip velocity

$$\dot{u} \approx 1.4 \frac{\tau}{\rho c} \frac{l_*}{\lambda},$$

(15)

during a typical earthquake. Using a shear stress $\tau \approx 10$ MPa, a density $\rho = 3000$ kg/m$^3$ and $c = 4$ km/s gives $\dot{u} \approx 1.2 (l_*/\lambda)$ m/s. Thus, a realistic slip velocity $\dot{u} \approx 1$ m/s requires that

$$l_* \approx \lambda.$$

(16)
Together with (11) and (9), this leads to

\[ Q \approx 3 \frac{\lambda}{l_e}. \]  

(17)

This last expression (17) is totally incompatible with the above condition \( Q \geq 10^3 \), as this would lead to \( l_e \approx \lambda/100 \) or smaller. This last condition is a physical impossibility: the elastic scattering length is always much larger than or at the extreme limit of the same order as the wavelength. The physical intuition is that a wave is defined over a length scale of the order of the wavelength (otherwise, there are no spatial oscillations) and the scattering process needs at least this scale to operate. The limit \( l_e \approx \lambda \) is attained only under exceptional circumstances leading to a novel phenomenon, called Anderson localization, in which the acoustic wave do not propagate anymore but oscillate locally. Extraordinary efficient scatterers are needed to reach this regime \([Sornette, 1989c]\). It is thus clear that the condition \( l_e \approx \lambda/100 \) is utterly unphysical.

If in contrast, we put \( Q \approx 10^3 \) in (10), we get \( l_\ast \approx 160 \lambda \), which from (15) leads to a maximum slip velocity \( \dot{u} \approx 7.5 \text{ mm/s} \), using the numerical example of Melosh [1996]. This slow sliding velocity is unrealistic for earthquakes.

### 4 Possible remedies

A first remedy is to relax the condition used by Melosh that the acoustic pressure needs to reach the overburden pressure in order to significantly affect the fault friction. We propose that only a small fraction \( \eta \) of it is enough to liquifiy the fault.

Indeed, it is well-established experimentally \([Biarez and Hicher, 1994]\) that the elastic modulii of granular media under large cyclic deformations are much lower than their static values. This effect occurs only for sufficiently large amplitudes of the cyclic deformation, typically for strains \( \epsilon_a \) above \( 10^{-4} \). At \( \epsilon_a = 10^{-3} \), the elastic modulii are halved and at \( \epsilon_a = 10^{-2} \), the elastic modulii are more than five times smaller than their static values. As a consequence, the strength of the granular medium is decreased in proportion. Melosh (private communication) also finds in laboratory experiments that a large decrease in elastic modulus is required to fit the measured flow rate of acoustically fluidized debris. This is consistent with flow in granular
material in a completely dilatent state and agrees with measurements reported in the literature for the elastic modulus of highly strained granular materials.

In absence of cohesion forces, the strength of a granular material is solely due to the effect of gravity weight that put grains in contact together and the resistance to shear is governed by Coulomb’s law according to which the shear stress at the threshold for sliding is proportional to the normal stress. The coefficient of proportionality is the friction coefficient. As we have discussed above, in presence of acoustic fluidization, the resistance to shear deformation is decreased as a consequence of the reduction of the effective elastic modulus. Correlatively, the threshold for sliding is also decreased (strength is often proportional to elastic modulus in brittle material). In order to capture these phenomena, we propose, what is maybe the simplest approximation, that the criterion for unlocking the fault is changed from Melosh’s criterion to the condition that the acoustic pressure needs only reach a fraction $\eta$ of the overburden pressure. Let us stress that the essence of our argument is that the acoustic energy is fed by the moving fault and thus the acoustic particle velocity adjust to the changing elastic modulus so that the overall acoustic energy is “externally” controlled by the rate of global elastic dissipation. The acoustic fluidization thus controls the sliding threshold rather than solely the acoustic particle velocity.

We need to estimate the strain created by the acoustic field. The acoustic pressure is related to the acoustic particle velocity $v$ by

$$p = \rho cv.$$  \hspace{1cm} (18)

Assuming

$$p = \eta \rho gh,$$  \hspace{1cm} (19)

this yields

$$v = \eta gh/c \approx 12 \text{ m/s}$$  \hspace{1cm} (20)

for $p \approx 200 \text{ MPa}$, a density $\rho = 3 \times 10^3 \text{ kg/m}^3$, a velocity $c = 4000 \text{ m/s}$ and $\eta = 0.1$. This corresponds to an acoustic wave displacement $u_a = v/2\pi f \approx 2 \times 10^{-3}$ m at a frequency $f \sim 10^3 \text{ Hz}$. The corresponding strain $u_a/w$ is $\sim 2 \times 10^{-3}$ for a gouge width $w$ of the order of one meter [Melosh, 1996] over which the intense shaking occurs. These estimations suggest that Melosh’s criterion that the acoustic stress fluctuations
must approach the overburden stress on the fault for acoustic fluidization to occur is too drastic and smaller shaking can reduce significantly the fault friction.

Persuing this reasoning, we see that the fundamental equation (6) in [Melosh, 1996] is changed into

\[
d\Psi^2 \frac{d^2 \Psi}{d\zeta^2} = \Psi - r\Sigma^2 \left[ \frac{1 - \text{erf}(\eta\sqrt{\Psi})}{1 + \text{erf}(\eta\sqrt{\Psi})} \right], \tag{21}
\]

where \( \eta = 1 \) recovers the case treated by Melosh. \( \Psi \) is the normalized acoustic energy, \( \zeta = z/l_\ast \), \( z \) is the coordinate perpendicular to the fault, the regeneration parameter is \( r = eQ/2 \) where \( e \) is the acoustic energy conversion efficiency, \( \Sigma = \tau/\rho gh \) and \( \text{erf}(x) \) is the error function. We see that a factor \( \eta < 1 \) implies a more effective generation of acoustic waves because the second source term of the r.h.s. is larger. But, since the bracket term saturates to one for large energies and/or small \( \eta \), this does not lead to a significantly larger slip velocity than found above. This remains a problem of the theory.

This problem might be alleviated by treating \( e \) self-consistently as a decreasing function of the friction coefficient, and thus of the acoustic energy density. The problem then becomes even more non-linear because it reflects in addition the dependence of the acoustic radiation efficiency of the granular gouge on the amplitude of the acoustic vibrations. In addition, the elastic modulus is also really nonlinear and it is only the tangent modulus that decreases close to the yield in the dilatent region, which suggests that the above linear formalism is not adequate and should be modified.

Further improvement could also take into account that the stochastic acoustic energy may deviate from a Rayleigh distribution [Ishimaru, 1978; Mirlin et al., 1998] when the medium is strong heterogeneous. This modifies the functional form of the term in bracket in eq.(21) and thus all numerical estimations.

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