

Fractal solids, product measures and fractional wave equations

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This paper builds on the recently begun extension of continuum thermomechanics to fractal porous media that are specified by a mass (or spatial) fractal dimension D , a surface fractal dimension d and a resolution length scale R . The focus is on pre-fractal media (i.e. those with lower and upper cut-offs) through a theory based on a dimensional regularization, in which D is also the order of fractional integrals employed to state global balance laws. In effect, the governing equations are cast in forms involving conventional (integer order) integrals, while the local forms are expressed through partial differential equations with derivatives of integer order but containing coefficients involving D , d and R . This procedure allows a specification of a geometry configuration of continua by ‘fractal metric’ coefficients, on which the continuum mechanics is subsequently constructed. While all the derived relations depend explicitly on D , d and R , upon setting $D = 3$ and $d = 2$, they reduce to conventional forms of governing equations for continuous media with Euclidean geometries. Whereas the original formulation was based on a Riesz measure—and thus more suited to isotropic media—the new model is based on a product measure, making it capable of grasping local fractal anisotropy. Finally, the one-, two- and three-dimensional wave equations are developed, showing that the continuum mechanics approach is consistent with that obtained via variational energy principles.

Keywords: fractal; fractional integrals; product measures; continuum mechanics; waves; fractal derivative

1. Introduction

Fractals date back to research by Hausdorff and Besicovich on monster sets over 100 years ago, and then to the seminal work of Mandelbrot (1982). He was then followed by physicists and mathematicians. The first category was primarily condensed matter physicists who focused on the effects of fractal geometries of materials on bulk responses, e.g. Feder (1988). A number of specialized models have also been developed to particular problems like wave scattering at fractals (Berry 1979), computational mechanics (Soare & Picu 2007), fracture mechanics (Chudnovsky & Kunin 1987; Balankin 1997; Carpinteri *et al.* 1999; Yavari *et al.* 2002*a,b*) or geomechanics (Dyskin 2004). While in recent years

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mathematicians began to look at partial differential equations—starting with Laplace’s or heat equation—on fractal (albeit non-random) sets (e.g. Kigami 2001; Strichartz 2006), an analogue of continuum physics and mechanics still needs to be developed. In particular, what is missing is a single unifying theoretical framework.

A new step in the direction of continuum physics and mechanics, relying on dimensional regularization, was taken by Tarasov (2005*a,b,c*). He developed continuum-type equations of conservation of mass, momentum and energy for fractal porous media, and on that basis studied several fluid mechanics and wave motion problems. In principle, one can then map a mechanics problem of a fractal (which is described by its mass (D) and surface (d) fractal dimensions plus the spatial resolution (R)) onto a problem in the Euclidean space in which this fractal is embedded, while having to deal with coefficients explicitly involving D , d and R . As it turns out, D is also the order of fractional integrals employed to state global balance laws. Clearly, this has very interesting ramifications for formulating continuum-type mechanics of fractal media, which needs to be further explored. The great promise stems from the fact that the conventional requirement of continuum mechanics, the separation of scales, does not hold, yet the partial differential equations (with derivatives of integer order) may still be employed.

This paper builds on Tarasov’s approach, which has already involved: an extension to continuum thermomechanics and thermoelasticity, a formulation of integral theorems, a generalization of extremum and variational principles and turbulence in fractal porous media (Ostoja-Starzewski 2007*a,b*, 2008, 2009, *in press*). We first specify the geometry configuration of continua via ‘fractal metric’ coefficients, and on that basis construct a continuum mechanics. Whereas the original formulation of Tarasov (and initially employed by us) was based on a Riesz measure—and thus more suited to isotropic fractals—the new model is based on a product measure, making it capable of grasping local fractal anisotropy. To verify our model, we require that the mechanical approach be consistent with that obtained from energy principles. To make the two approaches equivalent, we discuss the definition of fractal strain (with the fractal derivative developed) and verify product measures for anisotropic fractals. As an application, we study the one-dimensional, the two-dimensional anti-plane and the three-dimensional wave equations in general anisotropic fractal solids. While all the derived relations depend explicitly on D , d and R , upon setting $D = 3$ and $d = 2$, they reduce to conventional forms of governing equations for continuous media with Euclidean geometries.

2. Mass power law and fractional integrals

By a fractal solid, we understand a medium B having a fractal geometric structure. The mass of the medium m obeys a power law with respect to the lengthscale of measurement R (or resolution)

$$m(R) = kR^D, \quad D < 3, \quad (2.1)$$

where D is the fractal dimension of mass, and k is a proportionally constant. We note that, in practice, a fractional power law relation (2.1) is widely recognized and can be determined in experiments by a log–log plot of m and R (e.g.

Schroeder 1990), but the physical object is a pre-fractal, i.e. it has lower and upper cut-offs. Now, following Tarasov (2005a), the fractional integral is employed to represent mass in a three-dimensional region W

$$m(W) = \int_W \rho(\mathbf{r}) dV_D = \int_W \rho(\mathbf{r}) c_3(D, r) dV_3. \quad (2.2)$$

Here the first and the second equality involve fractional integrals and conventional integrals, respectively. The coefficient $c_3(D, r)$ provides a transformation between the two. Using Riesz fractional integrals, $c_3(D, r)$ is

$$c_3(D, r) = \frac{2^{3-D} \Gamma(3/2)}{\Gamma(D/2)} r^{D-3}, \quad r = \sqrt{\sum_{i=1}^3 (x_i)^2}. \quad (2.3)$$

Note that $c_3(D, r)$ above solely depends on the scalar distance r , which in turn confines the formulations to isotropic fractals. However, in general, the medium exhibits different fractal dimensions along different directions—it is anisotropic. A practical example of such a fractal anisotropy is given by Carpinteri *et al.* (1999, 2004), where a porous concrete structure is modelled by a Sierpinski carpet in the cross-section and a Cantor set along the longitudinal axis.

This consideration leads us to replace equation (2.1) by a more general power law relation with respect to each spatial coordinate

$$m(x_1, x_2, x_3) \sim x_1^{\alpha_1} x_2^{\alpha_2} x_3^{\alpha_3}. \quad (2.4)$$

In order to account for such anisotropies, the fractional integral representing mass distribution is specified via a product measure

$$m(x_1, x_2, x_3) = \iiint \rho(x_1, x_2, x_3) d\mu_1(x_1) d\mu_2(x_2) d\mu_3(x_3). \quad (2.5)$$

Here the length measurement $d\mu_k(x_k)$ in each coordinate is provided by

$$d\mu_k(x_k) = c_1^{(k)}(\alpha_k, x_k) dx_k, \quad k = 1, 2, 3. \quad (2.6)$$

The total fractal dimension of mass D then equals $\alpha_1 + \alpha_2 + \alpha_3$. It follows that the volume coefficient c_3 is given by

$$c_3 = c_1^{(1)} c_1^{(2)} c_1^{(3)} = \prod_{i=1}^3 c_1^{(i)}. \quad (2.7)$$

For the surface coefficient c_2 we typically consider a cubic volume element, each surface element of which is specified by the corresponding normal vector (along the axis i, j or k , see figure 1). Therefore, the coefficient $c_2^{(k)}$ associated with the surface $S_d^{(k)}$ is

$$c_2^{(k)} = c_1^{(i)} c_1^{(j)} = \frac{c_3}{c_1^{(k)}}, \quad i \neq j \text{ and } i, j \neq k. \quad (2.8)$$

The expressions of length coefficients $c_1^{(k)}$ depend on forms of specific fractional integrals. In the sequel, adopting a modified Riemann–Liouville fractional integral recently formulated by Jumarie (2005, 2008), we have

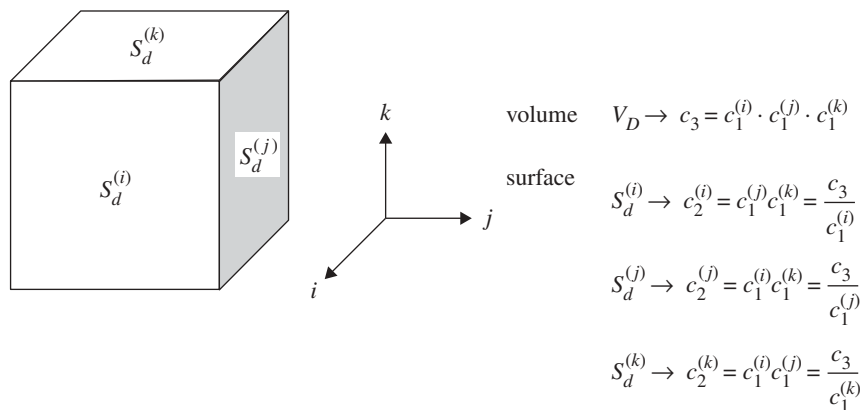


Figure 1. Constructing coefficients c_2 and c_3 via product measures.

$$c_1^{(k)} = \alpha_k (l_k - x_k)^{\alpha_k - 1}, \quad k = 1, 2, 3, \tag{2.9}$$

where l_k is the total length (integral interval) along axis x_k . Let us now examine it in two special cases.

- (i) *Uniform mass.* The mass is distributed uniformly in a cubic region W with a power law relation (2.4). Denoting the mass density by ρ_0 and the cubic length by l , we obtain

$$m(W) = \rho_0 l^{\alpha_1} l^{\alpha_2} l^{\alpha_3} = \rho_0 l^{\alpha_1 + \alpha_2 + \alpha_3} = \rho_0 l^D, \tag{2.10}$$

which is consistent with the mass power law (2.1).

- (ii) *Point mass.* The distribution of mass is concentrated at one point, so that the mass density is denoted by the Dirac function $\rho(x_1, x_2, x_3) = m_0 \delta(x_1) \delta(x_2) \delta(x_3)$. The fractional integral representing mass becomes

$$m(W) = \alpha_1 \alpha_2 \alpha_3 l^{\alpha_1 - 1} l^{\alpha_2 - 1} l^{\alpha_3 - 1} m_0 = \alpha_1 \alpha_2 \alpha_3 l^{D-3} m_0. \tag{2.11}$$

When $D \rightarrow 3$ ($\alpha_1, \alpha_2, \alpha_3 \rightarrow 1$), $m(W) \rightarrow m_0$ and the conventional concept of point mass is recovered (Temam & Miranville 2005). Note that using the Riesz fractional integral will always give zero (0^{D-3}) except when $D = 3$ (by letting $0^0 = 1$), which, on the other hand, shows a non-smooth transition of mass with respect to its fractal dimension. This also supports our choice of the non-Riesz-type expressions for $c_1^{(k)}$ in equation (2.9).

Note that the above expression $c_1^{(k)}$ shows that the length dimension, and hence the mass m , will involve an unusual physical dimension following from the fractional integral (2.5). This is understandable since in mathematics a fractal curve exhibits only a finite measure with respect to a fractal dimensional length unit (Mandelbrot 1982). Since, in practice, we prefer physical quantities to have usual dimensions, an alternative way to address this issue is to non-dimensionalize the coefficients $c_1^{(k)}$. Therefore, we suggest replacing $(l_k - x_k)$ by $(l_k - x_k)/l_0$ in equation (2.9) (l_0 is a characteristic scale, e.g. the mean pore size).

3. Fractional integral theorems and fractal derivatives

At this point, we recall two basic integral theorems extensively employed in continuum mechanics: Gauss's theorem, which relates a certain volume integral to the integral over its bounding surface, and the Reynolds transport theorem, concerning the rate of change of any volume integral for a continuous medium. In the following, we derive their fractional generalizations and, moreover, introduce a definition of fractal derivatives, which together provide a stepping-stone to construct a continuum mechanics in the setting of fractals.

The derivation of a fractional Gauss's theorem will be analogous to Tarasov's (2005*b*) dimensional regularization, albeit formulated in the framework of product measures discussed above. First, let us recall the surface integral in a fractal medium

$$S_d[\bar{f}] := \int_{S_d} \bar{f} \cdot \hat{n} \, dS_d = \int_{S_d} f_k n_k \, dS_d. \quad (3.1)$$

Here $\bar{f} = f_k \mathbf{e}_k$ is any vector field and $\hat{n} = n_k \mathbf{e}_k$ is the unit normal vector of the surface; the Einstein summation convention is assumed. In order to compute equation (3.1), we relate the integral element $\hat{n} \, dS_d$ to its conventional forms $\hat{n} \, dS_2$ via fractal surface coefficients $c_2^{(i)}, c_2^{(j)}, c_2^{(k)}$. Note that, by definition, any infinitesimal surface element dS_d in the integrand can be regarded as a plane (aligned in an arbitrary direction with a normal vector \hat{n}). Since the coefficients $c_2^{(i)}$'s are built on coordinate planes $Ox_j x_k$'s, we consider their projections onto each coordinate plane. The projected planes $n_i dS_d$ can then be specified by coefficients $c_2^{(i)}$'s, and this totally provides a representation of the integral element $\hat{n} \, dS_d$ (figure 2). Thus, we have

$$\int_{S_d} \bar{f} \cdot \hat{n} \, dS_d = \int_{S_2} f_k c_2^{(k)} n_k \, dS_2. \quad (3.2)$$

Now, following the conventional Gauss's theorem, we get

$$\int_{\partial W} f_k c_2^{(k)} n_k \, dS_2 = \int_W (f_k c_2^{(k)})_{,k} \, dV_3. \quad (3.3)$$

Note that from the expression (2.8) $c_2^{(k)}$ is independent of the variable x_k , so that we can write equation (3.3) as

$$\begin{aligned} \int_{\partial W} f_k n_k \, dS_d &= \int_W (f_k c_2^{(k)})_{,k} c_3^{-1} \, dV_D = \int_W f_{k,k} c_2^{(k)} c_3^{-1} \, dV_D \\ &= \int_W \frac{f_{k,k}}{c_1^{(k)}} \, dV_D := \int_W \nabla_k^D f_k \, dV_D. \end{aligned} \quad (3.4)$$

This equation is a fractional generalization of Gauss's theorem. Hereafter, we use the notation of fractal derivative, ∇_k^D , with respect to the coordinate x_k

$$\nabla_k^D := \frac{1}{c_1^{(k)}} \frac{\partial}{\partial x_k} (\cdot). \quad (3.5)$$

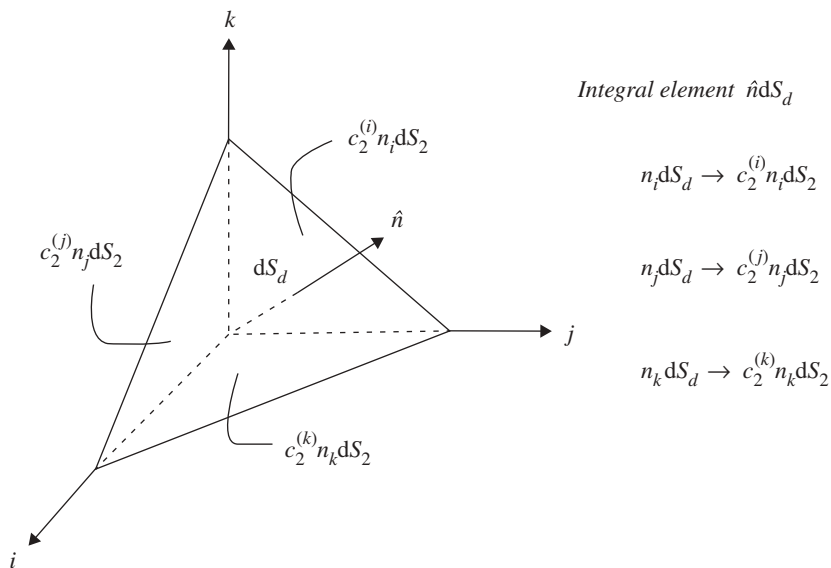


Figure 2. A representation of the fractional integral element $\hat{n} dS_d$ under product measures.

The definition of ∇_k^D is similar to Tarasov’s (2005b) $\nabla_k^D = c_3^{-1}(c_2 \cdot)_k$, albeit our form is simplified for product measures. We now examine three properties of the operator ∇_k^D .

- (i) It is the ‘inverse’ operator of fractional integrals. For any function $f(x)$, we have

$$\nabla_x^D \int f(x) d\mu^D(x) = \frac{1}{c_1(x)} \frac{d}{dx} \int f(x) c_1(x) dx = \frac{1}{c_1(x)} [f(x) c_1(x)] = f(x) \tag{3.6}$$

and

$$\int \nabla_x^D f(x) d\mu^D(x) = \int \left[\frac{1}{c_1(x)} \frac{df(x)}{dx} \right] c_1(x) dx = \int \frac{df(x)}{dx} dx = f(x). \tag{3.7}$$

For this reason, we call ∇_k^D a ‘fractal derivative’ (so as to distinguish it from the fractional derivatives already in existence).

- (ii) The rule of ‘term-by-term’ differentiation is satisfied

$$\nabla_k^D (AB) = \frac{1}{c_1^{(k)}} \frac{\partial}{\partial x_k} (AB) = \frac{1}{c_1^{(k)}} \frac{\partial(A)}{\partial x_k} B + \frac{1}{c_1^{(k)}} \frac{\partial(B)}{\partial x_k} A = B \nabla_k^D (A) + A \nabla_k^D (B), \tag{3.8}$$

whereby we note that this is invalid in Tarasov’s (2005b) notation.

- (iii) Its operation on any constant is zero

$$\nabla_k^D (C) = \frac{1}{c_1^{(k)}} \frac{\partial(C)}{\partial x_k} = 0. \tag{3.9}$$

Here we recall that the usual fractional derivative (Riemann–Liouville) of a constant does not equal zero—neither in fractional calculus (Oldham & Spanier 1974) nor in Tarasov’s (2005*b*) formulation.

As to the fractional generalization of Reynolds transport theorem, we follow the line of conventional continuum mechanics distinguishing between the reference and deformed configurations

$$\begin{aligned}
 \frac{d}{dt} \int_{W_t} P dV_D &= \frac{d}{dt} \int_{W_0} PJ dV_D^0 = \int_{W_0} \frac{d}{dt} (PJ) dV_D^0 = \int_{W_0} \left(\frac{d}{dt} P \cdot J + P \cdot \frac{d}{dt} J \right) dV_D^0 \\
 &= \int_{W_0} \left(\frac{d}{dt} P \cdot J + P \cdot v_{k,k} J \right) dV_D^0 = \int_{W_0} \left(\frac{d}{dt} P + P \cdot v_{k,k} \right) J dV_D^0 \\
 &= \int_{W_t} \left(\frac{d}{dt} P + P \cdot v_{k,k} \right) dV_D = \int_{W_t} \left(\frac{\partial}{\partial t} P + P_{,k} v_k + P \cdot v_{k,k} \right) dV_D \\
 &= \int_{W_t} \left(\frac{\partial}{\partial t} P + (P v_k)_{,k} \right) dV_D. \tag{3.10}
 \end{aligned}$$

Here P is any quantity accompanied by a moving material system W_t , $\mathbf{v} = v_k \mathbf{e}_k$ is the velocity field and J is the Jacobian of the transformation from the current configuration x_k , to the referential configuration X_K . Note that the result is identical to its conventional representation. The fractal material time derivative is thus the same

$$\left(\frac{d}{dt} \right)_D P = \frac{d}{dt} P = \frac{\partial}{\partial t} P + P_{,k} v_k, \tag{3.11}$$

whereby we note that the alternative form of the fractional Reynolds transport theorem which involves surface integrals is different from the conventional one and rather complicated. This is because the fractal volume coefficient c_3 depends on all coordinates x_k ’s (not like $c_2^{(k)}$, which is independent of x_k when deriving fractional Gauss’s theorem). Continuing with equation (3.10), the derivation follows as

$$\begin{aligned}
 \frac{d}{dt} \int_{W_t} P dV_D &= \int_{W_t} \left(\frac{\partial}{\partial t} P + (P v_k)_{,k} \right) dV_D = \int_{W_t} \frac{\partial}{\partial t} P dV_D + \int_{W_t} (P v_k)_{,k} c_3 dV_3 \\
 &= \int_{W_t} \frac{\partial}{\partial t} P dV_D + \int_{W_t} \left(\int (P v_k)_{,k} c_3 dx_k \right)_{,k} dV_3 = \int_{W_t} \frac{\partial}{\partial t} P dV_D \\
 &\quad + \int_{\partial W_t} \left(\int (P v_k)_{,k} c_3 dx_k \right) n_k dS_2 \\
 &= \int_{W_t} \frac{\partial}{\partial t} P dV_D + \int_{\partial W_t} \left(P v_k c_3 - \int P v_k c_{3,k} dx_k \right) (c_2^{(k)})^{-1} n_k dS_d \\
 &= \int_{W_t} \frac{\partial}{\partial t} P dV_D + \int_{\partial W_t} P c_1^{(k)} v_k n_k dS_d - \int_{\partial W_t} \left(\int P c_{1,k}^{(k)} v_k dx_k \right) n_k dS_d. \tag{3.12}
 \end{aligned}$$

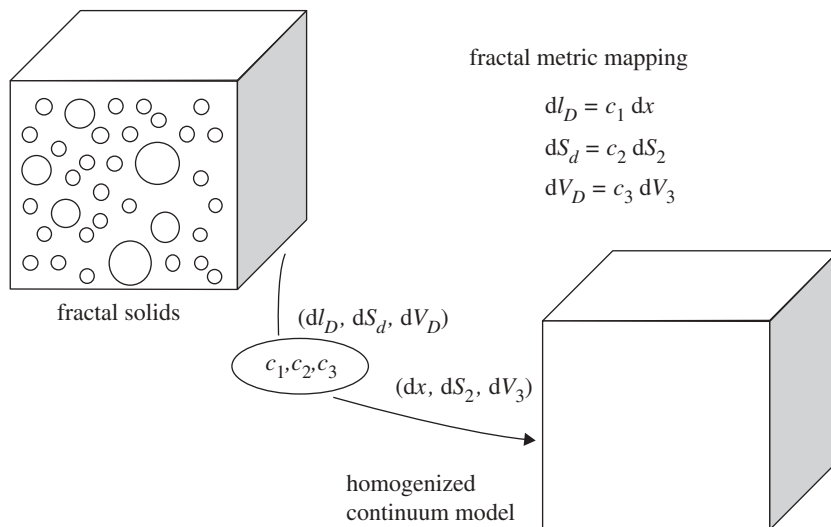


Figure 3. An illustration of the homogenization process from geometry configurations.

4. Continuum mechanics of fractal solids

In §§2 and 3, we have discussed fractional integrals under product measures and thereby generalized some basic integral theorems. Now, we proceed to develop a framework of continuum mechanics in fractal setting. We will formulate the field equations analogous to those in classical continuum mechanics but based on fractional integrals. Note that the notions of continuum mechanics rely on geometry configurations of the body. We shall first examine some physical concepts and definitions on account of the fractal geometry.

Let us recall the formula of fractal mass (equation (2.2)) which expresses the mass power law via fractional integrals. From a homogenization standpoint, this allows an interpretation of the fractal (intrinsically discontinuous) medium as a continuum and a ‘fractal metric’ embedded in the equivalent ‘homogenized’ continuum model, saying that

$$dl_D = c_1 dx, \quad dS_d = c_2 dS_2, \quad dV_D = c_3 dV_3. \tag{4.1}$$

Here dl_D, dS_d, dV_D represent the line, surface, volume element in the fractal body and dx, dS_2, dV_3 denote those in the homogenized model, see figure 3. The coefficients c_1, c_2, c_3 provide the relation between both sets.

The definitions of stress and strain must also be modified accordingly. The Cauchy stress is now specified to express the surface force F_k^S via fractional integrals

$$F_k^S = \int_{\partial W} \sigma_{kl} n_l dS_d = \int_{\partial W} \sigma_{kl} n_l c_2^{(l)} dS_2. \tag{4.2}$$

As to the configuration of strain, we recommend replacing all the spatial derivatives $\partial/\partial x_k$ with fractal derivatives ∇_k^D introduced in §3. This can be understood by observing from equation (4.1) that

$$\nabla_k^D = \frac{1}{c_1^{(k)}} \frac{\partial}{\partial x_k} = \frac{\partial}{\partial l_D^{(k)}}. \quad (4.3)$$

For small deformation, the expression of strain in fractal solids thus gives

$$\varepsilon_{ij} = \frac{1}{2} (\nabla_j^D u_i + \nabla_i^D u_j) = \frac{1}{2} \left(\frac{1}{c_1^{(j)}} u_{i,j} + \frac{1}{c_1^{(i)}} u_{j,i} \right). \quad (4.4)$$

Note that the stress–strain pairs must be conjugate from the viewpoint of energy. We shall examine the consistency of these definitions later when deriving wave equations in the next section. At this point, let us consider the balance law of linear momentum in fractal solids

$$\frac{d}{dt} \int_W \rho \mathbf{v} dV_D = \mathbf{F}^B + \mathbf{F}^S, \quad (4.5)$$

where $\mathbf{v} = v_k \mathbf{e}_k$ denotes the velocity vector, and \mathbf{F}^B , \mathbf{F}^S are the body and surface forces, respectively. Writing equation (4.5) in indicial notation and expressing forces in terms of fractional integrals yield

$$\frac{d}{dt} \int_W \rho v_k dV_D = \int_W f_k dV_D + \int_{\partial W} \sigma_{kl} n_l dS_d. \quad (4.6)$$

On observation of fractional Gauss's theorem (3.4) and Reynolds transport theorem (3.10), this gives

$$\int_W \rho \left(\frac{d}{dt} \right)_D v_k dV_D = \int_W (f_k + \nabla_l^D \sigma_{kl}) dV_D. \quad (4.7)$$

Here the operators of fractal derivative ∇_k^D and material derivative $(d/dt)_D$ are employed; recall equations (3.5) and (3.11). Note that the region W is arbitrary. On account of equation (4.7), we obtain the balance equation in local form

$$\rho \left(\frac{d}{dt} \right)_D v_k = f_k + \nabla_l^D \sigma_{kl}. \quad (4.8)$$

The specification of constitutive equations involves more arguments from physics, and we recommend keeping the relations of stress and strain while modifying their definitions to fractal setting. This is understood in that the fractal geometry solely influences our configurations of some physical quantities (like stress and strain) while it takes no effect on physical laws (like the conservation principles, and constitutive relations that are inherently due to material properties). We note that this justification is verified in Carpinteri & Pugno (2005) where scale effects of material strength and stress are discussed from the standpoint of fractal geometry and confirmed by experiments of both brittle and plastic materials.

Now, we consider a specific example: linear elastic solids with small deformation. The constitutive equations take the usual linear forms

$$\sigma_{ij} = \lambda \varepsilon_{kk} \delta_{ij} + 2\mu \varepsilon_{ij}, \quad (4.9)$$

where λ and μ are material parameters (Lame constants), while σ_{ij} and ε_{ij} are fractal stress and strain defined in equations (4.2) and (4.4), respectively.

Under small displacements, the linearization of equation (4.8) gives

$$\rho \frac{\partial^2 u_k}{\partial t^2} = f_k + \nabla_l^D \sigma_{kl}, \quad (4.10)$$

where $\mathbf{u} = u_k \mathbf{e}_k$ is the displacement field. Note that equations (4.4), (4.9) and (4.10) constitute a complete set of equations describing the problem (excluding boundary conditions).

5. Fractional wave equations

It is now possible to study wave motion in fractal solids based on the continuum-type equations derived in §4. As a starting point here, we will exclusively consider waves in linear elastic homogeneous fractal solids under small motions and zero external loads. Equations (4.4), (4.9) and (4.10) can thus jointly lead to wave equations. Note that variational principles provide an alternative approach to studying elastic problems. We shall therefore derive fractional wave equations via these two approaches and examine whether the results are consistent. This can be regarded as a verification of our entire formulation.

(a) *One-dimensional plane wave*

The one-dimensional plane wave motion involves one spatial variable only, say x_1 or just x . We will consider the derivations under mechanical and variational approaches, respectively. It is examined in the simplest case whether our definitions of fractal stress and strain in §4 are self-consistent.

(i) *Mechanical approach*

The balance of linear momentum reduces to

$$\rho \ddot{u} = c_1^{-1} \sigma_{,x}. \quad (5.1)$$

The constitutive equation becomes

$$\sigma = E \varepsilon, \quad (5.2)$$

where we recognize Young's modulus E . Substituting equation (5.2) into equation (5.1), we obtain

$$\rho \ddot{u} = E c_1^{-1} \varepsilon_{,x}. \quad (5.3)$$

Note that the strain ε is defined as a function of the displacement u (usually the derivative). The wave equation can then be derived from equation (5.3). Following

the conventional strain definition, $\varepsilon = u_{,x}$, which substituted into equation (5.3) gives

$$\rho \ddot{u} = E c_1^{-1} u_{,xx}. \quad (5.4a)$$

On the other hand, using our definition (4.4), simplified to one dimension, $\varepsilon = c_1^{-1} u_{,x}$, which yields

$$\rho \ddot{u} = E c_1^{-1} (c_1^{-1} u_{,x})_{,x}. \quad (5.4b)$$

(ii) *Variational approach*

In the variational approach, we consider the kinetic energy T and the strain energy U associated with the medium. The wave equation follows from Hamilton's principle that acts on its Lagrangian function $L = T - U$. First, the kinetic energy is

$$T = \frac{1}{2} \rho \int \dot{u}^2 dl_D = \frac{1}{2} \rho \int \dot{u}^2 c_1 dx, \quad (5.5)$$

while the strain energy is

$$U = \frac{1}{2} E \int \varepsilon^2 dl_D = \frac{1}{2} E \int \varepsilon^2 c_1 dx. \quad (5.6)$$

Employing the conventional definition of strain, equation (5.6) becomes

$$U = \frac{1}{2} E \int c_1 u_{,x}^2 dx, \quad (5.7a)$$

while using our fractal definition of strain gives

$$U = \frac{1}{2} E \int c_1^{-1} u_{,x}^2 dx. \quad (5.7b)$$

According to Hamilton's principle, $\delta \int L dt = \delta \int (T - U) dt = 0$, which implies the Euler-Lagrange equation

$$\frac{\partial}{\partial t} \left[\frac{\partial \ell}{\partial \dot{u}} \right] + \frac{\partial}{\partial x} \left[\frac{\partial \ell}{\partial u_{,x}} \right] - \frac{\partial \ell}{\partial u} = 0, \quad (5.8)$$

where ℓ denotes the Lagrangian density, defined by $L = \int \ell dx$. For the conventional definition of strain, ℓ has the form

$$\ell = \frac{1}{2} (\rho c_1 \dot{u}^2 - E c_1 u_{,x}^2), \quad (5.9a)$$

while the fractal definition of strain gives

$$\ell = \frac{1}{2} (\rho c_1 \dot{u}^2 - E c_1^{-1} u_{,x}^2). \quad (5.9b)$$

Substituting equation (5.9a) or (5.9b) into equation (5.8), we obtain, respectively, the wave equations

$$\rho c_1 \ddot{u} - E (c_1 u_{,x})_{,x} = 0, \quad (5.10a)$$

$$\rho c_1 \ddot{u} - E (c_1^{-1} u_{,x})_{,x} = 0. \quad (5.10b)$$

Comparing the results between the mechanical and variational approaches, we find that equation (5.4*b*) agrees with (5.10*b*), while equation (5.4*a*) contradicts (5.10*a*) (Tarasov 2005*b,c*). Thus, our definitions of fractal stress and strain are self-consistent.

An analogous result, also exhibiting self-consistency, has recently been obtained in elastodynamics of fractally structured Timoshenko beams (Ostoj-Starzewski & Li in press).

(*b*) *Two-dimensional anti-plane wave*

A two-dimensional anti-plane wave is described by a displacement field $u_3(x_1, x_2, t)$ (u_1 and u_2 vanishes). As before, we study it under mechanical and variational approaches. This then provides a simple case to examine our construction of the surface coefficient c_2 via product measures in §2.

For the local balance of linear momentum, only the one involving u_3 is of interest

$$\rho \ddot{u}_3 = \nabla_k^D \sigma_{3k,k}. \quad (5.11)$$

Note that the mean strain ε_{kk} is zero, and so the constitutive equations reduce to

$$\sigma_{ij} = 2\mu \varepsilon_{ij}. \quad (5.12)$$

The corresponding stress components in equation (5.11) follow from equations (5.12) and (4.4) as

$$\sigma_{3k} = \mu \nabla_k^D u_3 = \mu \frac{u_{3,k}}{c_1^{(k)}}. \quad (5.13)$$

Substituting equation (5.13) into equation (5.11), we obtain the wave equation

$$\rho \ddot{u}_3 = \mu \left[\frac{1}{c_1^{(1)}} \left(\frac{u_{3,1}}{c_1^{(1)}} \right)_{,1} + \frac{1}{c_1^{(2)}} \left(\frac{u_{3,2}}{c_1^{(2)}} \right)_{,2} \right]. \quad (5.14)$$

As to the variational approach, we consider the body with a unit length in x_3 . The kinetic energy thus gives (since only surface coefficient $c_2^{(3)}$ is involved, for simplicity we denote it as c_2)

$$T = \frac{1}{2} \rho \iint \dot{u}_3^2 dS_d = \frac{1}{2} \rho \iint \dot{u}_3^2 c_2 dS_2. \quad (5.15)$$

The strain energy is

$$U = \frac{1}{2} \iint \sigma_{ij} \varepsilon_{ij} dS_d = \frac{1}{2} \mu \iint (\nabla_k^D u_3) (\nabla_k^D u_3) c_2 dS_2. \quad (5.16)$$

The Lagrangian density has the form

$$\begin{aligned} \ell &= \frac{1}{2} \rho c_2 \dot{u}_3^2 - \frac{1}{2} \mu c_2 (\nabla_k^D u_3) (\nabla_k^D u_3) = \frac{1}{2} \rho c_1^{(1)} c_1^{(2)} \dot{u}_3^2 - \frac{1}{2} \mu c_1^{(1)} c_1^{(2)} \left[\left(\frac{u_{3,1}}{c_1^{(1)}} \right)^2 + \left(\frac{u_{3,2}}{c_1^{(2)}} \right)^2 \right] \\ &= \frac{1}{2} \rho c_1^{(1)} c_1^{(2)} \dot{u}_3^2 - \frac{1}{2} \mu \left[\frac{c_1^{(2)}}{c_1^{(1)}} u_{3,1}^2 + \frac{c_1^{(1)}}{c_1^{(2)}} u_{3,2}^2 \right]. \end{aligned} \quad (5.17)$$

Next, applying the Euler–Lagrange equation

$$\frac{\partial}{\partial t} \left[\frac{\partial \ell}{\partial \dot{u}_3} \right] + \sum_{k=1}^2 \frac{\partial}{\partial x_k} \left[\frac{\partial \ell}{\partial u_{3,k}} \right] - \frac{\partial \ell}{\partial u_3} = 0, \quad (5.18)$$

we derive a wave equation

$$\rho c_1^{(1)} c_1^{(2)} \ddot{u}_3 - \mu \left[c_1^{(2)} \left(\frac{u_{3,1}}{c_1^{(1)}} \right)_{,1} + c_1^{(1)} \left(\frac{u_{3,2}}{c_1^{(2)}} \right)_{,2} \right] = 0, \quad (5.19)$$

and note that it is equivalent to equation (5.14). This verifies our expression for c_2 via product measures in equation (2.8).

(c) Three-dimensional waves

We now proceed to discuss the most general case: three-dimensional waves that involve all spatial variables: x_1, x_2, x_3 . Similar to the above derivations, in the mechanical approach we eliminate the stresses by displacements via strain-displacement relations (4.4) and constitutive laws (4.9), and then arrive at the wave equation from equation (4.10). The results are a little more complicated and have the form

$$\rho \ddot{u}_i = \mu \nabla_j^D \nabla_j^D u_i + (\lambda + \mu) \nabla_i^D \nabla_j^D u_j,$$

i.e.

$$\rho \ddot{u}_i = \mu \frac{1}{c_1^{(j)}} \left(\frac{u_{i,j}}{c_1^{(j)}} \right)_{,j} + (\lambda + \mu) \frac{u_{j,j}}{c_1^{(i)} c_1^{(j)}}. \quad (5.20)$$

On the other hand, in the variational approach the Lagrange density follows as

$$\begin{aligned} \ell &= \frac{1}{2} \rho c_3 \dot{u}_i \dot{u}_i - \frac{1}{2} c_3 \sigma_{ij} \varepsilon_{ij} = \frac{1}{2} \rho c_3 \dot{u}_i \dot{u}_i - \left[\frac{1}{2} \lambda c_3 (\varepsilon_{kk})^2 + \mu c_3 \varepsilon_{ij} \varepsilon_{ij} \right] \\ &= \frac{1}{2} \rho c_3 \dot{u}_i \dot{u}_i - \frac{1}{2} c_3 \left[\lambda (\nabla_k^D u_k)^2 + \mu (\nabla_j^D u_i \nabla_j^D u_i + \nabla_j^D u_i \nabla_i^D u_j) \right] \\ &= \frac{1}{2} \rho c_3 \dot{u}_i \dot{u}_i - \frac{1}{2} c_3 \left[\lambda \left(\frac{u_{k,k}}{c_1^{(k)}} \right)^2 + \mu \left(\left(\frac{u_{i,j}}{c_1^{(j)}} \right) \left(\frac{u_{i,j}}{c_1^{(j)}} \right) + \left(\frac{u_{i,j}}{c_1^{(j)}} \right) \left(\frac{u_{j,i}}{c_1^{(i)}} \right) \right) \right]. \end{aligned} \quad (5.21)$$

The Euler–Lagrange equations are given by

$$\frac{\partial}{\partial t} \left[\frac{\partial \ell}{\partial \dot{u}_i} \right] + \sum_{j=1}^3 \frac{\partial}{\partial x_j} \left[\frac{\partial \ell}{\partial u_{i,j}} \right] - \frac{\partial \ell}{\partial u_i} = 0. \quad (5.22)$$

Substituting equation (5.21) into (5.22), we find

$$\rho c_3 \ddot{u}_i - \mu \left(\frac{c_3 u_{i,j}}{(c_1^{(j)})^2} \right)_{,j} - (\lambda + \mu) \left(\frac{c_3 u_{j,i}}{c_1^{(i)} c_1^{(j)}} \right)_{,j} = 0. \quad (5.23)$$

On account of the formulation of product measures, we have

$$\frac{c_3}{(c_1^{(j)})^2} = \frac{c_2^{(j)} c_1^{(j)}}{(c_1^{(j)})^2} = \frac{c_2^{(j)}}{c_1^{(j)}} \quad \text{and} \quad \frac{c_3}{c_1^{(i)} c_1^{(j)}} = c_1^{(k)}. \quad (5.24)$$

Note that $c_2^{(j)}$ and $c_1^{(k)}$ are both independent of x_j . Equation (5.23) thus becomes

$$\rho \ddot{u}_i - \mu \frac{c_2^{(j)}}{c_3} \left(\frac{u_{i,j}}{c_1^{(j)}} \right)_{,j} - (\lambda + \mu) \frac{c_1^{(k)} u_{j,ij}}{c_3} = 0. \quad (5.25)$$

From equations (2.7) and (2.8), it is clear that the results of the mechanical approach (5.20) are consistent with those of the variational approach (5.25) (and $u_{j,ji} = u_{j,ij}$). On the other hand, we note that in Tarasov's (2005*a,b*) expressions for c_1, \dots, c_3 —where Riesz fractional integrals were adopted—the forms of fractional wave equations are more complicated and they are not equivalent under these two approaches. This and other comments in this paper are not meant as a criticism of Tarasov's work as, indeed, we have been very much motivated by his research.

6. Conclusions

We presented a self-consistent approach to mechanics of continua with fractal geometries. The fractal characteristics are captured by fractional integrals following from the mass power law of the body. On this basis, a framework of continuum mechanics of fractal media is developed and a particular example concerning one-dimensional, two-dimensional anti-plane and three-dimensional waves in linear elastic fractal solids is studied, showing that the result of the mechanical approach is equivalent to that obtained by variational principles, thus verifying the self-consistency of our theoretical framework. Finally, our method suggests specific expressions for c_1, \dots, c_3 and, therefore, can be easily applied in practice when the power law of fractal mass is measured.

We note that the framework remains self-consistent whenever the expressions for c_1, \dots, c_3 are constructed via product measures and such a property is independent of a specific form of c_1 . This allows the modification of c_1 to possibly incorporate the asymptotic behaviours in lower and upper cut-offs of physical pre-fractals. Although our current focus is on fractal solids, the work can be potentially extended to fluid mechanics like the shear flow turbulence, e.g. regarding the anisotropic drag reducing flow (McComb & Chan 1979, 1981).

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CORRECTION

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Fractal solids, product measures and fractional wave equations

BY J. LI AND M. OSTOJA-STARZEWSKI

Equations (5.20)₂ and (5.25) are incorrect and should read as shown below. First, note that the summation convention involving indices in coefficients $c_1^{(j)}$, $c_2^{(k)}$ and c_3 is somewhat different from the usual one and should be carefully handled. Thus, in equation (5.20), corresponding to the first equation, the second equation should read

$$\rho \ddot{u}_i = \mu \frac{1}{c_1^{(j)}} \left(\frac{u_{i,j}}{c_1^{(j)}} \right)_j + (\lambda + \mu) \frac{1}{c_1^{(i)}} \left(\frac{u_{j,j}}{c_1^{(j)}} \right)_i. \quad (5.20)_2$$

Although $c_1^{(j)}$ is a function of x_j only, and independent of x_i ($i \neq j$), the term $u_{j,j}/c_1^{(j)}$ is subject to summation over $j = 1, 2, 3$ and equals $u_{1,1}/c_1^{(1)} + u_{2,2}/c_1^{(2)} + u_{3,3}/c_1^{(3)}$. Therefore, it is incorrect to pull $c_1^{(j)}$ out from the derivative. Similarly, in equation (5.23) the expression $(c_3 u_{j,i}/c_1^{(i)} c_1^{(j)})_j$ involves a summation over $j = 1, 2, 3$ so that, effectively, equation (5.25) should be disregarded and we obtain the same result as in equation (5.20)₂.

To recapitulate, the coefficients $c_1^{(i)}$, $c_1^{(j)}$, $c_1^{(k)}$ and $c_2^{(i)}$, $c_2^{(j)}$, $c_2^{(k)}$ are not involved in the summation convention. For example, equation (3.2) is written explicitly as

$$\int_{S_d} \bar{f} \cdot \hat{n} dS_d = \int_{S_2} f_k c_2^{(k)} n_k dS_2 = \sum_{k=1}^3 \int_{S_2} f_k c_2^{(k)} n_k dS_2 \quad (3.2)$$

owing to the repeated index in f_k and n_k . Also note that equation (3.5) does not involve a summation on k , although $c_1^{(k)}$ and x_k have a repeated index ‘ k ’

$$\nabla_k^D := \frac{1}{c_1^{(k)}} \frac{\partial}{\partial x_k} (\cdot) \neq \sum_{k=1}^3 \frac{1}{c_1^{(k)}} \frac{\partial}{\partial x_k} (\cdot). \quad (3.5)$$