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1 **Dynamic mass density of resonant metamaterials with homogeneous inclusions**

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Abstract

7 The occurrence of a negative dynamic mass density is a striking property of meta-
8 materials. It appears when an inner local resonance is present. Results coming from an
9 asymptotic theory are recalled briefly, showing the scaling of physical properties leading to
10 inner resonance in elastic composites containing homogeneous soft inclusions, with negli-
11 ble scattering of waves travelling through the matrix. This appears for a large contrast of
12 elastic properties between matrix and inclusion. The frequency-dependent dynamic mass
13 density depends on the resonance frequencies of the inner inclusions and on their related
14 participation factors. Having solved the dynamic elasticity problem, these physical quan-
15 tities are provided in the case of homogeneous cylindrical and spherical inclusions. It is
16 shown that numerous resonance frequencies do not contribute to the dynamic mass density
17 or have small participation factors, which simplifies significantly the physics involved in the
18 concerned inner resonance phenomena. Finally, non-dimensional resonance frequencies and
19 participation factors are given for both cases of inclusions as functions of the Poisson's ratio,
20 defining completely the dynamic mass density.

21 I. INTRODUCTION

22 Even if the existence of metamaterials has been predicted initially in the case of
23 electromagnetic properties^{41;36}, acoustic metamaterials have been studied extensively
24 during the last few decades, more specially since the production of experiments showing
25 band gaps when studying the propagation of elastic waves in phononic crystals⁴⁰ followed a
26 few years after by experiments of wave propagation through acoustic metamaterials²⁰ .
27 Acoustic metamaterials can have more or less complex internal structures as prove the
28 numerous papers in this field either studying properties of metamaterials from a theoretical
29 approach^{17;18;19;7;11} and/or showing experimentally the properties of these
30 materials^{20;38;28;46;33;44}. These materials can have "effective" negative dynamic density
31 and/or elastic moduli for some frequency ranges, leading to forbidden frequency bands. A
32 recent review paper on acoustic metamaterials²³ draws a large picture of this domain,
33 showing that the interesting properties of metamaterials are enhanced by inserting local
34 resonators. These resonators can be a more or less complex combination of plates, beams,
35 spring-mass systems, Helmholtz resonators or (homogeneous or heterogeneous) inclusions of
36 different shapes. This review paper reports also several recent studies on the conception of
37 materials having simultaneously negative effective mass density and elastic moduli^{16;14;18} .

38 In the following, the study will be restricted to elastic composites containing
39 inclusions of different shapes. Even so, the amount of literature devoted to this field is

40 impressive. It is known since a long time that every periodic composite containing
41 inclusions is characterized by forbidden frequency bands, its response to wave propagation
42 being studied usually by Bloch waves. However, these forbidden frequencies appear at high
43 frequencies, when the wavelength is of the order of the inclusion size,i.e. the material being
44 considered as a "phononic crystal". A few decades ago, it was recognized that the
45 frequency range of forbidden frequency bands is lowered by introducing inner resonators
46 and that the most simple way to induce inner resonance is to use matrix-inclusion
47 composites containing rigid inclusions surrounded by a soft coating (i.e. "composite
48 inclusions"), the whole being immersed within an elastic matrix. This led to the concept of
49 acoustic metamaterials. The composite inclusions act as "spring-mass" resonators, the
50 mass being provided by the inner inclusion and the spring by the surrounding coating. In
51 the case of a large contrast between elastic properties of coating and matrix, the low
52 resonance frequency of the spring-mass resonator can be reached practically without
53 scattering of the waves by the inclusions. It is the extreme case of "inner-resonance with
54 negligible scattering". Obviously, when the elastic properties of the coating become of the
55 same order as the one of the matrix, it is necessary to take into account the coupling of the
56 inner resonance with the scattering of waves by the inclusions. Several methods were used
57 to account for this coupling^{48;47;43;21;45;24}, the most currently used being the Coherent
58 Potential Approximation³⁷. However, in the case of a soft coating, the case of "inner

59 resonance with negligible scattering” is certainly an interesting comparison model, which
60 becomes fully valid for large contrasts of elastic properties of constituents.

61 Less work was devoted to homogeneous inclusions, and these works deal mainly with
62 scattering by inner inclusions. However, early works using homogenization theory based on
63 asymptotic expansion^{1;2} have shown that ”inner resonance with negligible scattering” can
64 occur also in composites containing homogeneous inclusions. This is the domain of
65 application of the present paper.

66 The physics of wave propagation through composite materials is complex and a strong
67 physical insight can be gained by using an asymptotic series expansion⁸ of the dynamic
68 solution for different assumptions of the ratios between physical parameters: elastic
69 coefficients and mass densities. First homogenization results showing the occurrence of
70 bandgaps in periodic elastic composites with ”inner resonance with negligible scattering”
71 were obtained in early works^{1;2}, but with the development of the ideas on metamaterials,
72 new results were obtained with the use of asymptotic expansions in dynamic
73 elasticity^{34;4;5;3;39;6;10;32;31;25;26}. More specifically, Auriault and Boutin³ have recently
74 extended the method of asymptotic expansion to materials containing composite inclusions
75 and characterized by ”inner resonance with negligible scattering”. It is worthwhile noticing
76 that, compared with results coming from other publications in this field, an important
77 aspect does appear in the papers written by Auriault and al. that are at the basis of this

78 paper^{1;2;3}: it is shown indeed in these papers that *the use of different scalings on physical*
79 *properties can lead to different macroscopic behaviours*. In the case of a strong contrast, the
80 method proves the occurrence of "inner resonance with negligible scattering" ; it provides
81 the structure of the effective constitutive equations and the local elasticity problems
82 (usually called "localization problems" or "cell problems") to be solved in order to obtain
83 the macroscopic physical properties, more particularly the dynamic mass density. It is
84 worth noticing that in absence of strong contrasts, the predicted dynamic behaviour can be
85 non-local in space^{12;42;29;30}, while the behaviour described in the case of "inner resonance
86 with negligible scattering" is local in space (while non-local in time). This point is of
87 importance for modelling wave propagation through a structure which would be made of
88 such a metamaterial.

89 In a recent paper⁹ the dynamic mass density has been obtained by solving the
90 localization problem described by the homogenization theory in the case of materials
91 displaying "inner resonance with negligible scattering" and containing spherical or
92 cylindrical composite inclusions that act as "spring-mass" resonators. Under suitable
93 conditions, such a dynamic mass density can be obtained by studying the static behaviour
94 (spring effect) of the composite inclusion. By comparison, in the case of homogeneous
95 inclusions described thereafter, the dynamic mass density involves the full dynamic
96 behaviour of the inclusions.

97 In section 2, the results of asymptotic expansions in the case of materials containing
98 homogeneous inclusions are synthesized and discussed, showing that suitable ratios of the
99 physical parameters can lead to inner resonance with negligible scattering. The method
100 specifies also the dynamic boundary value problem ("cell problem") to be solved in order
101 to provide the dynamic mass density. In section 3, the inner motion of inclusions made of
102 cylindrical fibers is studied: the eigenfrequencies are obtained and then participation
103 factors and dynamic mass densities of these resonators are given in a closed form. The case
104 of spherical inclusions is studied in section 4 and the full solution of the local elasticity
105 problem is given in this case, providing again the eigenfrequencies, participation factors
106 and dynamic mass densities. Finally, numerical applications are presented in section 5.

107 **II. DYNAMIC BEHAVIOUR OF ELASTIC COMPOSITES CONTAINING** 108 **HOMOGENEOUS INCLUSIONS**

109 This section recalls briefly the main results from homogenization using asymptotic
110 theory^{1;2;3} and the underlying assumptions that justify the method used in the following
111 sections to obtain the dynamic mass density. Let us consider a periodic elastic composite
112 material made up of a matrix containing homogeneous inclusions. This composite material
113 is defined by the geometry of the periodic cell containing the inclusions and by the physical
114 properties of the constituent materials. These constituents are assumed elastic and
115 isotropic. They are characterized by their mass densities $\rho^{(s)}$, their volume concentrations

116 $c^{(s)}$ and their Lamé elastic parameters $\lambda^{(s)}, \mu^{(s)}$. (s) corresponds to the matrix (m) or the
 117 inclusion (i). The effective behaviour of the composite depends strongly on the scaling
 118 parameters. In a first step, these scaling parameters are defined and next, the scaling
 119 assumptions leading to the occurrence of a negative mass density will be provided.

120 **A. Scaling parameters**

121 A first parameter is the geometric scaling ratio defined by $\epsilon = \frac{l}{L}$, where l is the size of
 122 the periodic cell and L is the order of magnitude of the wavelength within the matrix in a
 123 chosen frequency range. The case of interest is when the size of the periodic cell is small
 124 compared with L and therefore $\epsilon \ll 1$.

125 Other scaling parameters contain the various physical parameters that characterize
 126 the constituents of the composite.

127 They comprise:

- 128 - the ratio between the mass densities of the constituents $\frac{\rho^{(i)}}{\rho^{(m)}}$.
- 129 - the ratio between the orders of magnitude of the elastic coefficients of the
 130 different constituents.

131 The order of magnitude of the elasticity coefficients of a given material will be defined
 132 by the value of $a^{(s)} = \lambda^{(s)} + 2\mu^{(s)}$ (i.e. a norm of the elasticity tensor) . The scaling ratio
 133 related to elastic coefficients will be defined by $\frac{\lambda^{(i)}+2\mu^{(i)}}{\lambda^{(m)}+2\mu^{(m)}} = \frac{a^{(i)}}{a^{(m)}}$:

134 **B. Scaling assumptions leading to the occurrence of a negative dynamic mass**

135 **density**

136 Using asymptotic expansion along the scaling parameter ϵ , Auriault and Bonnet¹
 137 (see also^{2;3}) studied the effect of different values of the previously defined scaling ratios on
 138 the overall behaviour of the composite. These authors showed that the resonance of the
 139 inner inclusions can be obtained when the inclusions are very soft compared with the
 140 matrix, the densities being of the same order of magnitude. More precisely, the inner
 141 resonance is obtained when the scaling ratios meet the following conditions: $\frac{a^{(i)}}{a^{(m)}} = O(\epsilon^2)$
 142 for the elastic parameters and $\frac{\rho^{(i)}}{\rho^{(m)}} = O(1)$ for the mass densities.

143 **C. The dynamic mass density**

Under the previously defined conditions, the composite has an overall effective
 behaviour for harmonic time excitation at radial frequency ω that is characterized by the
 dynamic equation

$$\frac{\partial(a_{ijkl}^{(eff)} \varepsilon_{kl})}{\partial x_j} + \omega^2 \rho_{ij}^{(eff)} u_j = 0 \quad (1)$$

144 where $a_{ijkl}^{(eff)}$ are the "effective" elastic coefficients of the composite, ε_{kl} are the components
 145 of the overall strain tensor, u_j are the components of the displacement and $\rho_{ij}^{(eff)}$ are the
 146 components of a frequency-dependent dynamic mass density. With a periodic array of
 147 inclusions, the effective elastic behaviour is not isotropic, contrarily to the behaviour of the
 148 constituents. Due to the scaling assumptions, the inclusions are very soft compared with

149 the matrix. As a consequence, the effective elastic coefficients $a_{ijkl}^{(eff)}$ can be computed as for
 150 a matrix that contains voids in place of inclusions by usual means to obtain effective
 151 properties of heterogeneous elastic materials.

Concerning the components $\rho_{ij}^{(eff)}$ of the mass density, they must be considered as the
 ones of a tensor of second rank. In the following, it will be assumed that the inclusions
 have three orthogonal symmetry planes. In this case, the dynamic mass density is diagonal
 for a coordinate system whose axes are parallel to the symmetry planes, with:

$$\rho_{ii}^{(eff)}(\omega) = \langle \rho \rangle + c^{(i)} \rho^{(i)} h_{ii}(\omega) \quad (2)$$

152 where Einstein's convention summation must not be applied to the repeated index i in
 153 $\rho_{ii}^{(eff)}$, h_{ii} . $\langle \rho \rangle$ is the volume average of the density within the periodic cell and $c^{(i)}$ is the
 154 volume concentration of inclusions. h_{ii} are the components of a diagonal second order
 155 tensor.

The homogenization method provides the localization problems ("cell problems")
 allowing to compute the effective elastic properties and the tensor of components h_{ij} , as
 recalled in Appendix I. The asymptotic process proves that, at the first order, due to the
 large wavelength within the matrix, the fluctuation of displacement of the matrix $\vec{u}^{(m)}$
 within the periodic cell is of the order of ϵ (while the fluctuations of strain field are of the
 same order as the macroscopic strain). It means that the displacement and acceleration of
 the matrix are nearly constant inside a periodic cell. As a consequence, the localization

problem obtained by homogenization theory recalled in Appendix I corresponds to a dynamic localization problem within the inclusion moved by a uniform acceleration induced by the matrix. It can be expressed by using the relative displacement

$\vec{w} = \vec{u}^{(i)} - \vec{u}^{(m)}$. This relative displacement is null over the boundary of the inclusion and is the solution of

$$\mu^{(i)} \Delta \vec{w} + (\lambda^{(i)} + \mu^{(i)}) \overrightarrow{\text{grad}} (\text{div}(\vec{w})) + \rho^{(i)}(\omega)^2 \vec{w} = -\rho^{(i)}(\omega)^2 \vec{u}^{(m)} \quad (3)$$

156 It means that the motion within a given inclusion, in a local reference frame moving with
157 the matrix, is impulsed by the inertial acceleration $-\rho^{(i)}(\omega)^2 \vec{u}^{(m)}$.

The solution of this problem can be expressed by using the eigenfrequencies ω^p and eigenmodes \vec{u}^p of the inclusions *for fixed boundaries in the reference frame moving with the matrix*, i.e. the solutions of

$$\mu^{(i)} \Delta \vec{u}^p + (\lambda^{(i)} + \mu^{(i)}) \overrightarrow{\text{grad}} (\text{div}(\vec{u}^p)) + \rho^{(i)}(\omega^p)^2 \vec{u}^p = 0 \quad (4)$$

158 with $\vec{u}^p = 0$ on the boundary of the inclusion.

Finally, for an acceleration of the matrix given by $\gamma \vec{e}_j$ where \vec{e}_j is a unit vector along one of the axes, the related component $h_{jj}(\omega)$ is given by:

$$h_{jj}(\omega) = \sum_{p=1}^{\infty} K_{jj}^p \cdot \frac{1}{\left(\frac{\omega^p}{\omega}\right)^2 - 1} \quad (5)$$

where the participation factors K_{jj}^p do not depend on ω and are given by

$$K_{jj}^p = \frac{(\langle \vec{u}^p \rangle^{(i)} \cdot \vec{e}_j)^2}{\langle \|\vec{u}^p\|^2 \rangle^{(i)}}, \quad (6)$$

159 the volume average $\langle f \rangle^{(i)}$ being computed over the volume of the inclusion. Due to the
 160 dependence on ω , h_{jj} tends to minus infinity just above each resonance frequency. As a
 161 consequence, the related component of the overall dynamic mass density $\rho_{jj}^{(eff)}$ tends also
 162 to minus infinity and remains negative until its negative frequency dependent part is
 163 equilibrated by the static overall mass density $\langle \rho \rangle$ in equation (4).

164 It is worthwhile noticing that, due to the condition at the boundary of the inclusion
 165 (induced obviously by the assumption of "negligible scattering"), the contributions of
 166 different separate inclusions within a periodic cell are completely independent and can be
 167 summed in order to compute the overall dynamic mass density for any distribution of
 168 inclusions.

169 Finally, the eigenfrequencies and eigenmodes (ω^p, \vec{u}^p) must be obtained for each
 170 direction of \vec{e}_j by solving a "localized eigenmodes problem" (4). Next, the averages in (6)
 171 are computed to produce the three components of the participation factors. In the following
 172 sections, the localization problems will be solved for cylindrical and spherical inclusions.

173 **D. Physical discussion of the results**

174 The results obtained from asymptotic expansion may be surprising by some aspects.
 175 Indeed, these results show that the inclusion behaves as if the surrounding material moves
 176 uniformly at its boundary. In addition, the field inside the inclusion does not depend on
 177 the wavelength within the surrounding material and is therefore not affected by the

178 scattering of these waves by the inclusions.

179 Obviously, the displacement within the matrix is not exactly uniform, due to the
180 varying strain around the inclusion. However, the ratio of strain fields between matrix and
181 inclusion is very large, due to the elasticity contrast, and as a consequence, the fluctuation
182 of displacement around the inclusion is very small compared with the displacement inside
183 the inclusion and may be considered as negligible. The absence of scattering comes
184 naturally from the initial assumption of a large ratio between the wavelength related to the
185 propagation within the matrix and the size of inclusions. Finally, the results coming from
186 the asymptotic theory appear as physically sound. Indeed, numerous devices entering the
187 composition of metamaterials use inner resonators whose resonance appears at low
188 frequencies, below frequencies inducing a scattering by the cells containing the resonators
189 (Helmholtz resonators, spring-mass resonators,...) .

190 Another point of discussion is the fact that the dynamic density tends to infinity at
191 the resonance frequencies. Obviously, this result comes from the assumption of perfectly
192 elastic materials. For real materials, the physical damping must be taken into account.
193 This can be effected by assuming the inclusion material as viscoelastic.

194 In this case, the viscoelastic behaviour of the inclusion material can be taken into
195 account by using complex and frequency dependent elastic moduli. As a result of the use of
196 complex moduli, the orthogonality of modes is no more ensured. Due to this loss of

197 orthogonality, the method presented previously cannot be used as such in the case of
198 viscoelasticity and will not be studied thereafter. For a slight damping, an asymptotic
199 procedure can again be used in order to provide the viscoelastic solution³. It leads to a
200 finite value of the dynamic density and to a small shift of the resonance frequency.

201 **III. DYNAMIC MASS DENSITY FOR CYLINDRICAL INCLUSIONS**

202 Let us consider a composite containing long parallel cylindrical inclusions, i.e. long
203 cylindrical fibers of radius a . As seen in section 2, the dynamic contribution to the
204 dynamic mass density depends only on the properties of the inclusions. Therefore, in the
205 following, the notation $\lambda^{(i)} = \lambda, \mu^{(i)} = \mu, \rho^{(i)} = \rho$ will be used. In a first step, the general
206 dynamic displacement within the fiber is given by using separate variables. Next, it will be
207 shown that numerous contributions to this general field do not provide contributions to the
208 participation factors. Finally, the expression of the dynamic mass density will be given.

209 **A. Dynamic displacement field within a cylindrical domain**

210 Different solutions exist for the dynamic displacement field within a cylindrical
211 domain^{22;15} using separate variables. The expression below is given in Eringen and
212 Suhubi¹⁵(eq.8.9.17-8.9.18) where it is shown that the displacement field is the

213 superposition for integer values of n of the following partial displacements:

$$\begin{aligned}
 u_r &= \frac{1}{r} [A_1 U_1(\alpha r) \cos n\theta e^{\pm i\gamma_p z} + B_1 U_2(\beta r) \sin n\theta e^{\pm i\gamma_q z} + C_1 U_3(\beta r) \cos n\theta e^{\pm i\gamma_q z}] \\
 u_\theta &= \frac{1}{r} [A_1 V_1(\alpha r) \sin n\theta e^{\pm i\gamma_p z} + B_1 V_2(\beta r) \cos n\theta e^{\pm i\gamma_q z} + C_1 V_3(\beta r) \sin n\theta e^{\pm i\gamma_q z}] \\
 u_z &= A_1 W_1(\alpha r) \cos n\theta e^{\pm i\gamma_p z} + C_1 W_3(\beta r) \cos n\theta e^{\pm i\gamma_q z}
 \end{aligned}$$

214 for $n = 0..∞$, where $\alpha = \frac{\omega}{c_p}$, $\beta = \frac{\omega}{c_s}$ are the wave numbers related to the celerities of

215 compressional and shear waves given by $c_p = \sqrt{(\lambda + 2\mu)/\rho}$, $c_s = \sqrt{\mu/\rho}$, with an

216 "alternative expression" obtained by replacing the *cos* terms by *sin* ones (and reversely)

217 and the constants A_1, B_1, C_1 by another set of constants A_2, B_2, C_2 .

218 $U_1, U_2, U_3, V_1, V_2, V_3, W_1, W_3$ are functions of r which are expressed by using Bessel
 219 functions.

220 As shown previously, the displacement field is sought for an acceleration parallel to
 221 one of the base vectors \vec{e}_j . This displacement field has the symmetry with respect to the

222 planes containing two coordinate axes. As a consequence, the symmetry induces that

223 $\gamma_p = \gamma_q = 0$. From the expression of the radial components, it comes that U_3, V_3 and W_1

224 are null. As another consequence of the symmetry, the constants B_1, A_2, C_2 are null.

225 Finally, replacing A_1, B_2, C_1 by A, B, C , it comes that the displacement field is given by:

$$\begin{aligned} u_r &= \frac{1}{r}[AU_1(\alpha r) \cos(n\theta) + BU_2(\beta r) \cos(n\theta)] = u_r^0 \cos(n\theta) \\ u_\theta &= \frac{1}{r}[AV_1(\alpha r) \sin(n\theta) + BV_2(\beta r) \sin(n\theta)] = u_\theta^0 \sin(n\theta) \\ u_z &= CW_3(\beta r) \cos(n\theta) \end{aligned}$$

226 The radial dependence is described by the following functions:

$$\begin{aligned} U_1(\alpha r) &= \alpha r J_{n-1}(\alpha r) - n J_n(\alpha r) \\ U_2(\beta r) &= n J_n(\beta r) \\ V_1(\alpha r) &= -n J_n(\alpha r) \\ V_2(\beta r) &= -\beta r J_{n-1}(\beta r) + n J_n(\beta r) \\ W_3(\beta r) &= \beta^2 J_n(\beta r) \end{aligned}$$

227 where J_n is the n^{th} order Bessel function of first kind. It can be seen that this displacement
 228 field contains two independent fields: a first one with components (u_r, u_θ) and a second one
 229 corresponding to u_z . The first one corresponds to a transversal displacement and the
 230 second one to a longitudinal displacement.

231 At this stage, the expression of the displacement field contains the components related
 232 to any value of n . However, it will be shown now that only a few of them contribute to the
 233 participation factors.

234 **B. Selection of the modes contributing to the dynamic mass density**

235 As seen in section 2, the participation factors are functions of the average of the
 236 projection of the displacement field along the direction of the unit vector \vec{e}_j . In a first
 237 step, this average projection is studied, because it is important to select only the modes
 238 which contribute to the dynamic mass density.

239 **1. Transversal motion**

240 It is assumed that the acceleration is transversal, i.e. \vec{e}_j is perpendicular to the axis of the
 241 cylinder. Due to the symmetry, the direction of \vec{e}_j can be chosen as \vec{e}_x corresponding to
 242 $\theta = 0$. The average of the projection of the displacement field over the section is given by
 243 I_t/S , where $S = \pi a^2$ and I_t is the integral of the projection of the displacement field onto
 244 $\vec{e}_x = \cos \theta \vec{e}_r - \sin \theta \vec{e}_\theta$ given by $I_t = \int \vec{u} \cdot \vec{e}_x r dr d\theta$, where
 245 $\vec{u} \cdot \vec{e}_x = u_r^0 \cos n\theta \cos \theta - u_\theta^0 \sin n\theta \sin \theta$. The integral over θ is null except for $n = 1$. It
 246 implies that only the contributions corresponding to $n = 1$ must be taken into account in
 247 the computation of the dynamic mass density. Finally, the transversal displacement field is
 248 given by:

$$u_r = \frac{1}{r} [AU_1(\alpha r) \cos \theta + BU_2(\beta r) \cos \theta] = u_r^0 \cos \theta$$

$$u_\theta = \frac{1}{r} [AV_1(\alpha r) \sin \theta + BV_2(\beta r) \sin \theta] = u_\theta^0 \sin \theta$$

249 with

$$\begin{aligned} u_r^0 &= \frac{1}{r}[AU_1(\alpha r) + BU_2(\beta r)] \\ u_\theta^0 &= \frac{1}{r}[AV_1(\alpha r) + BV_2(\beta r)] \end{aligned}$$

250 and

$$\begin{aligned} U_1(\alpha r) &= \alpha r J_0(\alpha r) - J_1(\alpha r) = r \frac{dJ_1(\alpha r)}{dr} \\ U_2(\beta r) &= J_1(\beta r) \\ V_1(\alpha r) &= -J_1(\alpha r) \\ V_2(\beta r) &= -\beta r J_0(\beta r) + J_1(\beta r) = -r \frac{dJ_1(\beta r)}{dr} \end{aligned}$$

251 **2. Longitudinal motion**

252 The longitudinal motion is obtained for $\vec{e}_j = \vec{e}_z$. In this case, the integral of the projection
 253 of the displacement field is given by $I_l = \int \vec{u} \cdot \vec{e}_z r dr d\theta = \int u_z r dr d\theta$. It involves the
 254 integral over θ of $\cos n\theta$ which is null except for $n = 0$. The longitudinal motion is therefore
 255 given by $u_z = CW_3(\beta r) = DJ_0(\beta r)$, where D is an undetermined constant.

256 **C. Computation of the eigenfrequencies**

257 **1. Transversal motion**

258 The eigenfrequencies correspond to the displacement fields which comply with the
 259 condition of null displacement over the boundary of the cylinder ($r = a$) and therefore to
 260 the linear system

$$AU_1(\alpha a) + BU_2(\beta a) = 0$$

$$AV_1(\alpha a) + BV_2(\beta a) = 0$$

This homogeneous system has a solution only when its determinant is null and finally the
 eigenfrequencies equation is given by $U_1(\alpha a)V_2(\beta a) - U_2(\beta a)V_1(\alpha a) = 0$. Using the
 expressions of functions U_i, V_i , this relation can be written by using only the
 non-dimensional frequency $\omega^* = \frac{\omega a}{c_s}$ as:

$$-k_c \omega^* J_0(k_c \omega^*) J_0(\omega^*) + k_c J_0(k_c \omega^*) J_1(\omega^*) + J_0(\omega^*) J_1(k_c \omega^*) = 0 \quad (7)$$

261 where $k_c = c_s/c_p$ is the ratio between celerities of shear and compressional waves, which
 262 depends only on the Poisson's ratio ν and is given by $k_c = \sqrt{\frac{1-2\nu}{2(1-\nu)}}$.

263 The solutions of this equation in ω^* are the non-dimensional eigenfrequencies
 264 ω^{p*} ($p = 1.. \infty$) related to the transversal modes. Having computed these eigenfrequencies,
 265 the eigenmodes are obtained by noticing that A and B must comply with the homogeneous
 266 equations and therefore $B = q^p \cdot A$ with $q^p = -U_1(k_c \omega^{p*})/U_2(\omega^{p*})$. These eigenmodes are
 267 thus defined up to the multiplying constant A .

268 The expression of the displacement field is no more valid for the case of an
 269 incompressible material, which corresponds to $k_c = 0$. However, the limit of the equation
 270 for eigenfrequencies when k_c tends to 0 can be obtained and leads to $J_2(\omega^*) = 0$. Therefore,
 271 the eigenfrequencies tend to the zeros of Bessel function J_2 when k_c tends to 0 (and also ν
 272 tends to 0.5).

273 **2. Longitudinal motion**

274 In this case, the non-dimensional eigenfrequencies are the solutions of $J_0(\omega^{p*}) = 0$. The
 275 non-dimensional eigenfrequencies ω^{p*} are therefore the zeros of Bessel function J_0 .

276 **D. Computation of the participation factors**

The participation factors K_{jj}^p are given by averages of the projection of the displacement field and of its quadratic norm. These averages involve the integrals of the projection $I_j = \int \vec{u} \cdot \vec{e}_j dS$ and of the square of norm of the displacement field $M_j = \int || \vec{u} ||^2 dS$. Using these quantities, K_p is given by:

$$K_{jj}^p = \frac{I_j^2}{S.M_j} \quad (8)$$

277 where S is the cross section area of the fiber. The integrals I_j and M_j are given below in a
 278 closed form.

279 **1. Transversal motion**

The integrals involved in the transversal motion are given by:

$$I_x = aA\pi[J_1(k_c\omega^{p*}) + q^p J_1(\omega^{p*})] \quad (9)$$

and

$$M_x = \pi A^2[G(k_c\omega^{p*}) + q^{p2}G(\omega^{p*}) + 2q^p J_1(k_c\omega^{p*})J_1(\omega^{p*})] \quad (10)$$

where

$$G(x) = \frac{1}{2}[(x^2 - 2)J_1^2(x) + x^2 J_0^2(x)] \quad (11)$$

A further simplification can be effected on the participation factors by taking into account the equation for eigenfrequencies (10). Finally, the participation factors are given by:

$$K_{xx}^p = \frac{2\alpha^*\beta^* J_2(\alpha^*)J_2(\beta^*)}{J_1(\alpha^*)J_1(\beta^*)(\alpha^{*2} + \beta^{*2}) - \alpha^*\beta^* [J_1(\beta^*)\beta^* J_0(\alpha^*) + J_1(\alpha^*)\alpha^* J_0(\beta^*)]} \quad (12)$$

280 where $\alpha^* = k_c\omega^{p*}$ and $\beta^* = \omega^{p*}$.

281 It is noteworthy that, in the case of incompressibility, the participation factors tend to
 282 0, because their numerators contain $J_2(\beta^*)$ which tends to zero with k_c as observed
 283 previously.

284 **2. Longitudinal motion**

In the same way, the participation factors for the longitudinal motion are given by:

$$I_z = \frac{2\pi Da^2}{\beta^2} J_1(\omega^{p*}) \quad (13)$$

$$M_z = \pi D^2 a^2 [J_0^2(\omega^{p*}) + J_1^2(\omega^{p*})] \quad (14)$$

$$K_{zz}^p = \frac{4J_1^2(\omega^{p*})}{\omega^{p*2} [J_0^2(\omega^{p*}) + J_1^2(\omega^{p*})]} \quad (15)$$

285 Taking into account the properties of eigenfrequencies leads finally to $K_{zz}^p = 4/\beta^{*2}$.

286 E. Synthesis of the results for cylindrical inclusions

287 The dynamic mass density (4) for cylindrical inclusions of radius a is characterized by
288 two sets of eigenfrequencies.

-The set of eigenfrequencies for transversal motion that are solutions of

$$-k_c \omega^* J_0(k_c \omega^*) J_0(\omega^*) + k_c J_0(k_c \omega^*) J_1(\omega^*) + J_0(\omega^*) J_1(k_c \omega^*) = 0 \quad (16)$$

289 where ω^* is the non dimensional frequency given by $\omega^* = \frac{\omega a}{c_s}$, k_c is given by $k_c = \sqrt{\frac{1-2\nu}{2(1-\nu)}}$

290 and J_0, J_1 are Bessel functions.

291 -The set of eigenfrequencies for longitudinal motion, solutions of $J_0(\omega^{p*}) = 0$.

292 The participation factors $K_{xx}^p = K_{yy}^p, K_{zz}^p$ for the transversal motion which enter the
293 expression of the dynamic density through (4), (6) and (7) are given

294

- for transversal motion by:

$$K_{xx}^p = \frac{2k_c J_2(k_c \omega^{p*}) J_2(\omega^{p*})}{J_1(k_c \omega^{p*}) J_1(\omega^{p*}) (1 + k_c^2) - k_c \omega^{p*} [J_1(\omega^{p*}) J_0(k_c \omega^{p*}) + k_c J_1(k_c \omega^{p*}) J_0(\omega^{p*})]} \quad (17)$$

295 -for longitudinal motion by $K_{zz}^p = 4/(\omega^{p*})^2$.

296 IV. DYNAMIC MASS DENSITY FOR SPHERICAL INCLUSIONS

297 We consider now a composite containing spherical inclusions having a radius a .
 298 Similarly to the previous section, it is necessary to compute the eigenmodes related to null
 299 displacement at the boundary of the inclusions. Numerous papers dealt with
 300 eigenfrequencies in the case of free boundary, but the case of fixed boundary is less
 301 frequently seen. This problem was studied early by Debye¹³ for the estimation of the
 302 phonon contribution to the specific heat of solids. Some eigenfrequencies for the fixed
 303 boundary problem were obtained also by Schafbuch et al.³⁵ However, to our knowledge, the
 304 computation of the participation factors and of the dynamic mass density has never been
 305 produced.

306 Taking into account the symmetry of the problem, the components of the dynamic
 307 mass density are the same for any orientation of the acceleration of the solid. Therefore,
 308 the participation factors K_{ij}^p are the components of an isotropic tensor with $K_{ij}^p = k^p \delta_{ij}$,
 309 where δ_{ij} is the Kronecker symbol.

310 A. Displacement field within a sphere

The displacement field within a sphere is given in classical books. The solution
 proposed in Eringen and Suhubi¹⁵(equations 8.13.14-8.13.15) for an harmonic motion leads
 to the following expression of the partial components of the displacement field in spherical

coordinates:

$$\begin{aligned}
 u_r &= \frac{1}{r} [AU_1(\alpha r) + CU_3(\beta r)] P_n^m(\cos \theta) \exp i(m\varphi) \\
 u_\theta &= \frac{1}{r} \left[[AV_1(\alpha r) + CV_3(\beta r)] [n \cot \theta \cdot P_n^m(\cos \theta) - \frac{n+m}{\sin \theta} P_{n-1}^m(\cos \theta)] \right. \\
 &\quad \left. + BV_2(\beta r) \frac{im}{\sin \theta} P_n^m(\cos \theta) \right] \exp i(m\varphi) \\
 u_\varphi &= \frac{1}{r} \left[[AV_1(\alpha r) + CV_3(\beta r)] \cdot \frac{im}{\sin \theta} P_n^m(\cos \theta) \right. \\
 &\quad \left. - rBV_2(\beta r) [n \cot \theta \cdot P_n^m(\cos \theta) - \frac{n+m}{\sin \theta} P_{n-1}^m(\cos \theta)] \right] \exp i(m\varphi)
 \end{aligned} \tag{18}$$

311 where A, B, C are constants, P_n^m are associated Legendre polynomials and the radial
 312 dependence is given by

$$\begin{aligned}
 U_1(\alpha r) &= nj_n(\alpha r) - \alpha r j_{n+1}(\alpha r) \\
 U_3(\beta r) &= n(n+1)j_n(\beta r) \\
 V_1(\alpha r) &= j_n(\alpha r) \\
 V_2(\beta r) &= j_n(\beta r) \\
 V_3(\beta r) &= (n+1)j_n(\beta r) - \beta r j_{n+1}(\beta r)
 \end{aligned} \tag{19}$$

313 These functions are expressed using spherical Bessel functions²⁷ j_n given by
 314 $j_n(kr) = \sqrt{\pi/2kr} J_{n+1/2}(kr)$. The computation of these spherical Bessel functions can be
 315 significantly alleviated, because these functions can be obtained by using trigonometric
 316 functions as: $j_1(z) = \sin z/z^2 - \cos z/z$ and $j_2(z) = (3/z^3 - 1/z)/\sin z - 3 \cos z/z^2$. The
 317 dynamic mass density does not depend on the direction of the matrix acceleration. It is

318 convenient to choose this direction along axis z . As for the case of cylindrical inclusions,
 319 the integral of $\vec{u} \cdot \vec{e}_z$ over the volume of the sphere must be different from zero for the
 320 components which contribute to the dynamic density. The computation of these integrals
 321 has been performed, leading to the result that the only contributions of the partial
 322 solutions correspond to $m = 0, n = 1$. All other components do not contribute to the
 323 dynamic mass density.

Finally, the dynamic displacement field is given by:

$$u_r = \frac{1}{r} [AU_1(\alpha r) + CU_3(\beta r)] \cos \theta$$

$$u_\theta = \frac{-1}{r} [AV_1(\alpha r) + CV_3(\beta r)] [\sin \theta]$$

324 with:

$$U_1(\alpha r) = j_1(\alpha r) - \alpha r j_2(\alpha r)$$

$$U_3(\beta r) = 2j_1(\beta r)$$

$$V_1(\alpha r) = j_1(\alpha r)$$

$$V_3(\beta r) = 2j_1(\beta r) - \beta r j_2(\beta r)$$

325 It is noteworthy that the structure of the displacement field presents strong similarities
 326 with the transversal displacement field within the cylinder, but involves spherical Bessel
 327 functions instead of usual Bessel functions.

328 **B. Resonance frequencies**

The non-dimensional resonance frequencies are obtained again by the determinant of the homogeneous system corresponding to a null displacement at $r = a$. It leads to the equation in $\omega^* = \frac{\omega a}{c_s}$:

$$k_c \omega^* j_2(k_c \omega^*) j_2(\omega^*) - j_1(k_c \omega^*) j_2(\omega^*) - 2k_c j_1(\omega^*) j_2(k_c \omega^*) = 0 \quad (20)$$

329 Having solved this equation and obtained the eigenfrequencies ω^{p*} , the boundary
330 condition leads also to the ratio $q^p = \frac{C}{A}$ for each eigenfrequency, with

$$\begin{aligned} q^p &= \frac{C}{A} = -\frac{U_1(k_c \omega^{p*})}{U_3(\omega^{p*})} = -\frac{V_1(k_c \omega^{p*})}{V_3(\omega^{p*})} \\ &= \frac{k_c \omega^{p*} j_2(k_c \omega^{p*}) - j_1(k_c \omega^{p*})}{2j_1(\omega^{p*})} \end{aligned}$$

331 As for the case of cylindrical inclusions, it is useful to study the case of incompressibility
332 when k_c tends to 0. In this case, the eigenfrequencies equation becomes $j_2(\omega^*) = 0$ and the
333 eigenfrequencies are the zeros of spherical Bessel function j_2 .

334 **C. Participation factors**

The participation factors are computed from the integrals $I = \int \vec{u} \cdot \vec{e}_z dV$ and
 $M = \int || \vec{u} ||^2 dV$. These quantities are given by:

$$I = \frac{4\pi A a^2}{3} [j_1(k_c \omega^{p*}) + 2q^p j_1(\omega^{p*})]$$

and

$$M = \int \|\mathbf{u}_p\|^2 = \frac{4\pi a A^2}{3} \{Q_{11}(k_c \omega^{p*}) + 4q^p j_1(k_c \omega^{p*}) j_1(\omega^{p*}) + (q^p)^2 Q_{33}(\omega^{p*})\}$$

335 where

$$Q_{11}(x) = j_1^2 - \frac{5}{2} x j_1 j_2 + \frac{x^2}{2} (j_1^2 + j_2^2)$$

$$Q_{33}(x) = 4j_1^2 - 5x j_1 j_2 + x^2 (j_1^2 + j_2^2)$$

336 with $j_1 = j_1(x)$, $j_2 = j_2(x)$.

As previously, the participation factors are finally given by introducing the averages

$I/V, M/V$ over the volume $V = \frac{4}{3}\pi a^3$ of the sphere by $k^p = I^2/V M$ or:

$$k^p = \frac{[j_1(k_c \omega^{p*}) + 2q^p j_1(\omega^{p*})]^2}{\{Q_{11}(k_c \omega^{p*}) + 4q^p j_1(k_c \omega^{p*}) j_1(\omega^{p*}) + q^{p2} Q_{33}(\omega^{p*})\}} \quad (21)$$

Introducing the values of the factors q^p and taking into account the equation for eigenfrequencies leads finally to:

$$k^p = \frac{2\alpha^* \beta^* j_2(\alpha^*) j_2(\beta^*)}{\alpha^* \beta^* (\beta^* j_1(\beta^*) j_2(\alpha^*) + \alpha^* j_1(\alpha^*) j_2(\beta^*)) - j_1(\alpha^*) j_1(\beta^*) (2\alpha^{*2} + \beta^{*2})} \quad (22)$$

337 When k_c tends to zero, the participation factors tend also to 0, because their numerators

338 contain $j_2(\beta^*) = j_2(\omega^{p*})$ which tends to 0 with k_c , as observed previously.

339 **D. Synthesis of the results for spherical inclusions**

The dynamic mass density (4) for spherical inclusions of radius a is characterized by the eigenfrequencies that are solutions of:

$$k_c \omega^* j_2(k_c \omega^*) j_2(\omega^*) - j_1(k_c \omega^*) j_2(\omega^*) - 2k_c j_1(\omega^*) j_2(k_c \omega^*) = 0 \quad (23)$$

340 where j_1, j_2 are spherical Bessel functions.

The participation factors which enter the expression of the dynamic density through (6) and (7) are given for each solution ω^{p*} of the previous equation by

$K_{xx}^p = K_{yy}^p = K_{zz}^p = k^p$, where k^p is given by:

$$k^p = \frac{2k_c j_2(k_c \omega^{p*}) j_2(\omega^{p*})}{k_c \omega^{p*} (j_1(\omega^{p*}) j_2(k_c \omega^{p*}) + k_c j_1(k_c \omega^{p*}) j_2(\omega^{p*})) - j_1(k_c \omega^{p*}) j_1(\omega^{p*}) (1 + 2k_c^2)} \quad (24)$$

341 V. APPLICATIONS

342 A. Cylindrical inclusions

343 As seen in sections 3 and 4 , the resonance frequencies and participation factors
 344 depend only on the non-dimensional frequency $\omega^* = \omega.a/c_s$. However, the coefficients
 345 appearing in the equation giving ω^* for the transversal motion depend on the ratio
 346 $k_c = c_s/c_p$, which itself depends only on Poisson's ratio ν .

347 Figure 1 displays the non-dimensional resonance frequencies for the transversal
 348 motion related to the first six modes as a function of the Poisson's ratio. The part near
 349 $\nu = 0.5$ has been enhanced on the right part of the figure. The resonance frequencies

350 increase moderately with ν for any mode, except when the Poisson's ratio is nearing
351 $\nu = 0.5$ where higher modes display a significant increase near this value of ν . In all cases,
352 the value of the eigenfrequency for $\nu = 0.5$ has been computed by using the zeros of J_2 ,
353 confirming the limit of the eigenfrequencies when ν tends to 0.5. All curves display a
354 plateau near $\nu = 0.5$, which is very short for the sixth resonance frequency, but is clear on
355 the enhanced figure at the right.

356 The participation factors related to these first six modes have been displayed on
357 Figure 2. The left part corresponds to all values of Poisson's ratios and the right part to ν
358 between 0.45 and 0.5. It shows that the first mode contains the main participation to the
359 dynamic mass density with a participation factor being near to 0.7 for small values of the
360 Poisson's ratio. However, when the Poisson's ratio increases, the first participation factor
361 decreases, the highest participation factor being successively the one related to increasing
362 ranks of modes, as show the values corresponding to ν nearing 0.5.

363 All participation factors tend to 0 for ν reaching 0.5, as expected from the expression
364 of the participation factors. These results show therefore that the use of nearly
365 incompressible materials leads to a prediction of the dynamic mass density which is very
366 sensitive to the value of the Poisson's ratio.

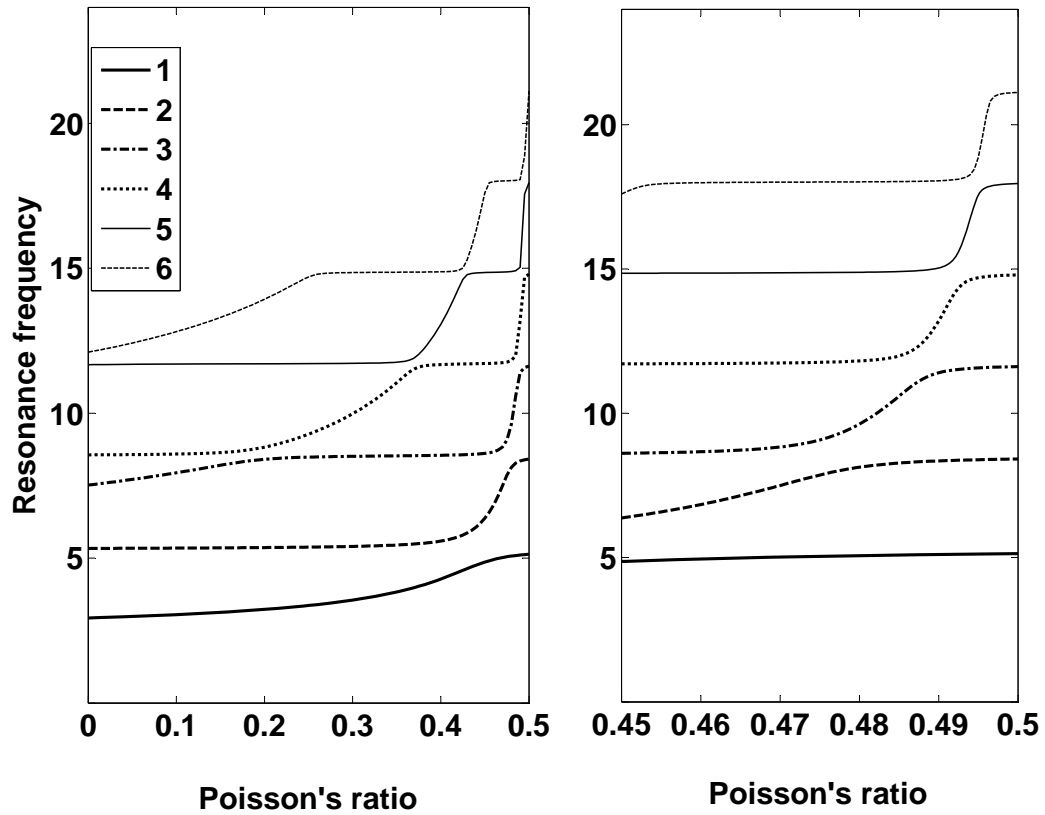


Figure 1: Cylindrical inclusions. First six non-dimensional resonance frequencies $\omega^{p*} = \frac{\omega^p a}{c_s}$, $p = 1..6$ for the transversal motion as functions of the Poisson's ratio ν . Right: enlargement of the range of ν between 0.45 and 0.5.

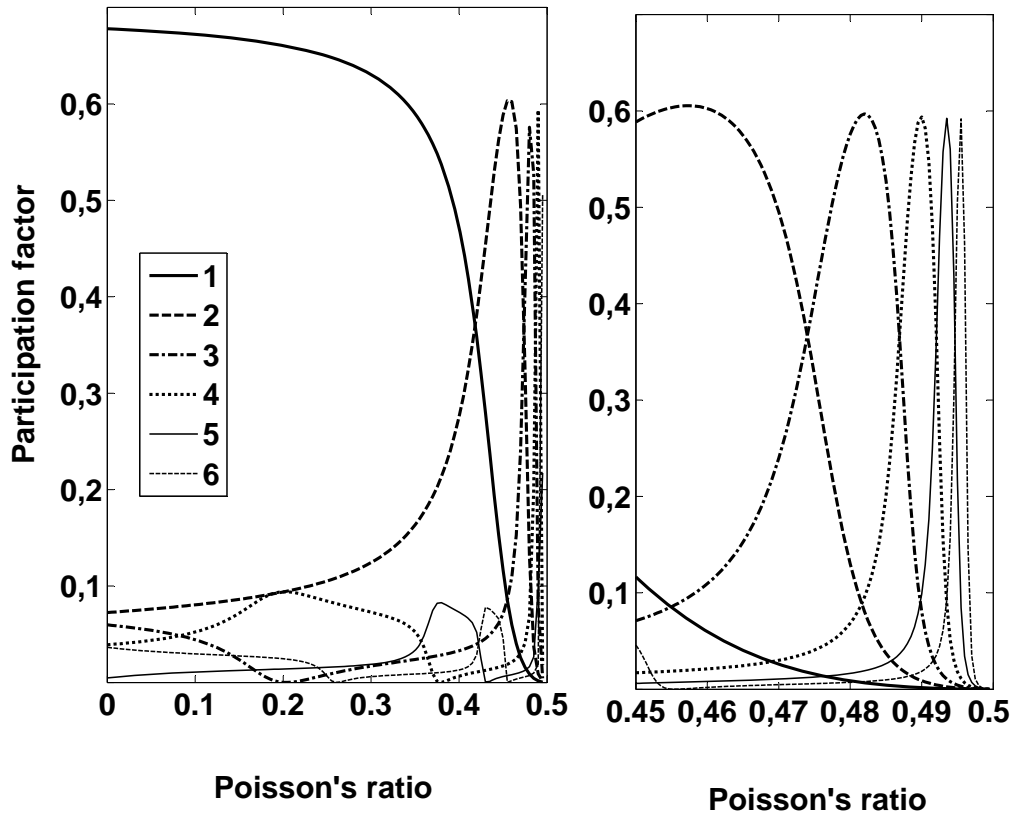


Figure 2: Cylindrical inclusions. First six participation factors for the transversal motion as functions of the Poisson's ratio ν . Right : enlargement of the range of ν between 0.45 and 0.5.

368 The nondimensional eigenfrequencies for spherical inclusions are reported in Figure 3
 369 and show values of eigenfrequencies that are higher than the ones for cylindrical inclusions
 370 with a trend being overall very similar to the one observed in the case of cylindrical
 371 inclusions.

372 The values of the related participation factors displayed on Figure 4 are also very
 373 similar to the ones observed in the case of cylindrical inclusions, the value for the first
 374 mode being however lower than for cylindrical inclusions, at low values of ν . The results
 375 show also a trend similar when the Poisson's ratio is nearing 0.5, the highest participation
 376 factors being related to higher modes. It is noteworthy that the participation factors are
 377 closed form expressions of the eigenfrequencies. However, the eigenfrequencies are more
 378 difficult to obtain. So, in order to make easier a further use of this work, precise values of
 379 nondimensional eigenfrequencies have been reported in Appendix II.

380 An example of dynamic mass density has been reported in Figure 5 for the case of a
 381 composite containing a matrix made of alumina (Young modulus $E_a = 350MPa$,
 382 $\nu_a = 0.25$, $\rho_a = 3950kgm^3$) and spherical inclusions made of polystyren
 383 ($E_p = 2MPa$, $\nu_p = 0.11$, $\rho_p = 900kg/m^3$) with a concentration of 0.5 and a radius of 2mm.
 384 The frequency range has been chosen to display the dynamic mass density ρ^{eff} around the
 385 first two eigenfrequencies, showing the ranges where ρ^{eff} is negative (highlighted by a bold
 386 line along the axis $\rho^{eff} = 0$). The frequency range of practical interest where ρ^{eff} is

387 negative could be enlarged by using reinforced polystyren containing inner inclusions with
388 a material of higher density, which would increase significantly the density of the inclusion
389 material and as a consequence the dynamic part of the density.

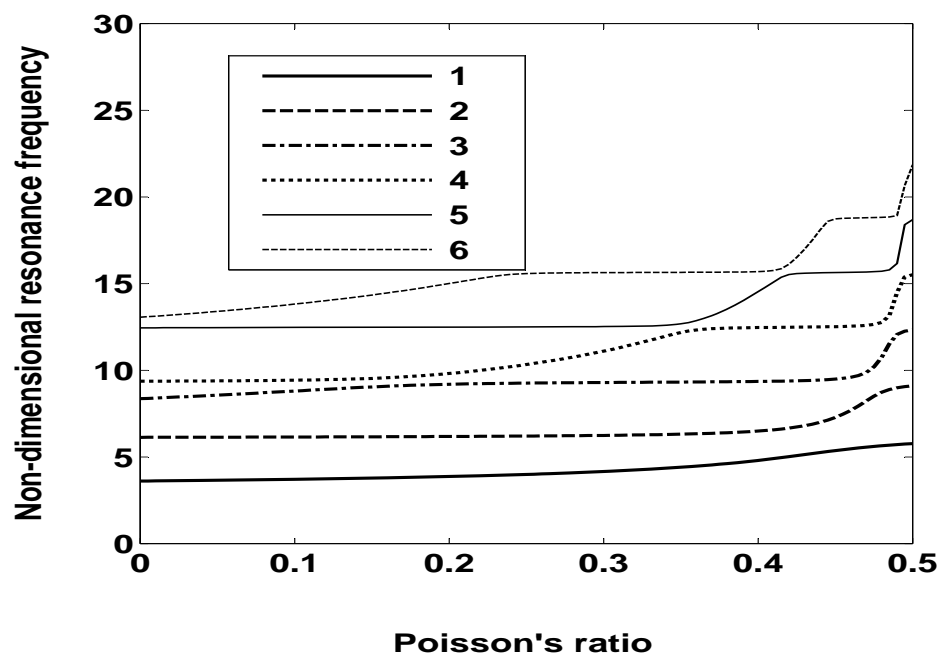


Figure 3: Spherical inclusions. First six non-dimensional resonance frequencies as functions of the Poisson's ratio ν

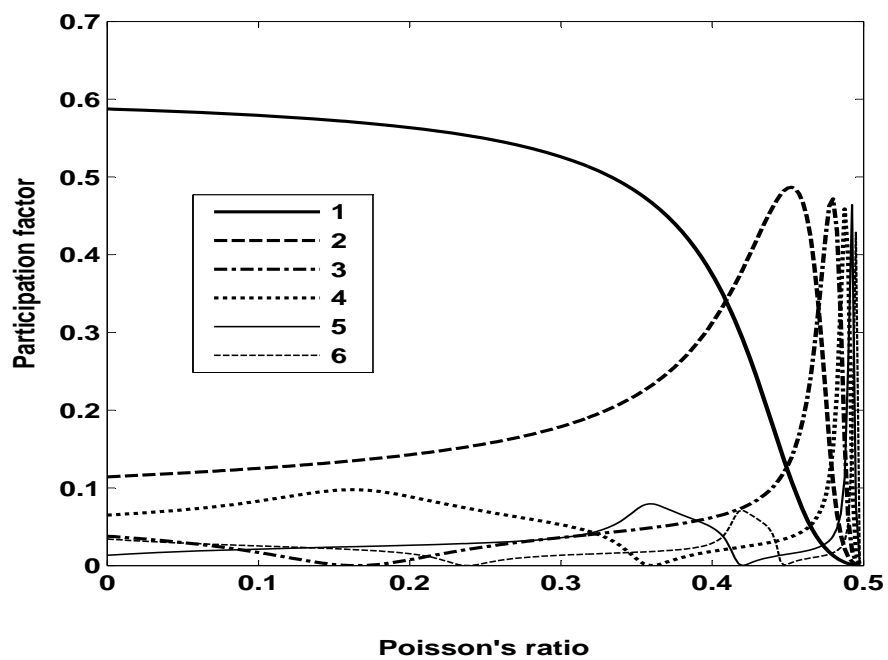


Figure 4: Spherical inclusions. First six participation factors as functions of the Poisson's ratio ν

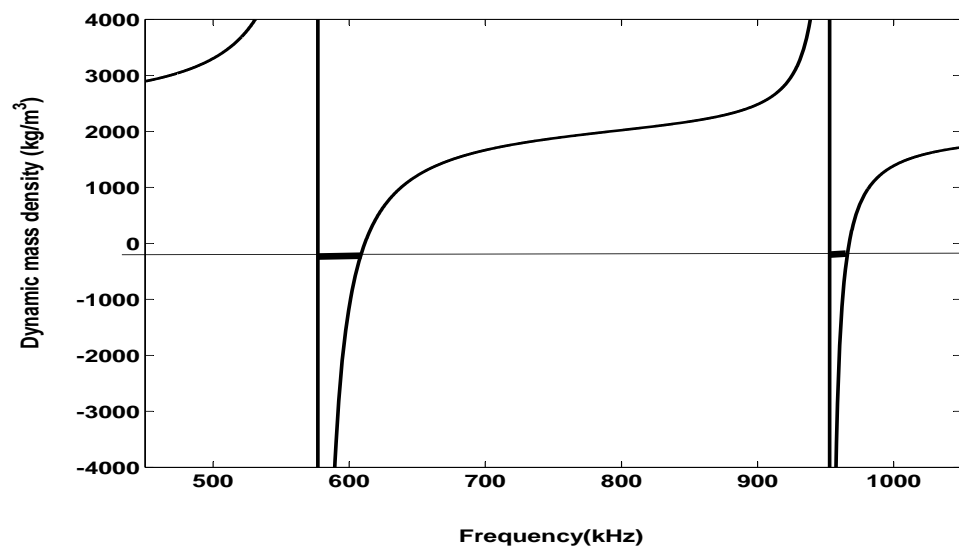


Figure 5: Dynamic mass density as a function of the frequency for a material with an alumina matrix containing spherical inclusions of polystyrene

390 **VI. CONCLUSION**

391 In this paper, the basic results obtained from the homogenization theory using
392 asymptotic expansions have been summarized for composite elastic metamaterials with
393 homogeneous inclusions displaying "inner resonance with negligible scattering". These
394 results show, under specific conditions corresponding to very soft inclusions with a
395 relatively high density, that the dynamic behaviour of the composite displays a
396 frequency-dependent and tensorial dynamic mass density. This dynamic mass density can
397 be computed knowing the resonance frequencies of the inclusions *for the condition of null*
398 *value at the surface of the inclusion of the relative displacement between matrix and*
399 *inclusion*. This greatly simplifies the computation of the dynamic mass density. In
400 addition, it shows that all resonators are independent.

401 The resonance frequencies of these resonators have been computed successfully in the
402 case of long cylindrical fibers or spherical inclusions. Having obtained these resonance
403 frequencies, the participation factors related to the different modes are given by simple
404 closed-form expressions. The results show that the participation factors decrease with
405 increasing ranks of resonance frequencies at moderate values of the Poisson's ratio. In this
406 case, the first modes constitute the main contribution to the dynamic mass density. When
407 the material is nearly incompressible, the contribution of modes related to higher
408 frequencies become more important and the dynamic mass density is very sensitive to the

409 exact value of the Poisson's ratio of the inclusions.

410 It can be noticed that the dynamic mass density tends to infinity at the resonance

411 frequencies, because the results have been obtained for a perfectly elastic material.

412 However, these results can be extended to viscoelastic materials in order to account for

413 physical damping, leading to a finite dynamic mass density by using an asymptotic solution

414 for a slight damping³. Such an extension will be the subject of a further paper.

415 **APPENDICES**

416 **I. THE CELL PROBLEM WITHIN THE INCLUSION**

417 The cell problem over the inclusion can be obtained^{1;2} by using a homogenization
 418 method based on double scale asymptotic expansion⁸, whose main aspects are recalled
 419 thereafter.

420 The microscopic structure is assumed periodic with a periodic cell Ω , splitted into Ω_m
 421 and Ω_i for matrix and inclusion, separated by the interface Γ between matrix and
 422 inclusion. One considers that the ordering is such that Ω_m is connected over all the
 423 domain, while Ω_i is made of separate inclusion domains.

424 The double scale method can be described as follows: the position of a point at the
 425 macroscopic scale is determined by using the so-called "slow position vector" \vec{x} , while the
 426 position of a point at the microscopic scale (i.e. inside the periodic cell) is determined by
 427 the "fast position vector" $\vec{y} = \vec{x}/\epsilon$, leading to an amplification of the fluctuation of physical
 428 variables when looking at the microscopic scale. The components of the displacement field
 429 are assumed to be functions of fast and slow variables by $\vec{u} = \vec{u}(\vec{x}, \vec{y})$ and their dependence
 430 in \vec{y} is periodic. Elasticity parameters and densities are also periodic functions of \vec{y} , by
 431 construction of the composite material.

432 The displacement field is searched by using an asymptotic expansion of the form:

433 $\vec{u}(\vec{x}, \vec{y}) = \vec{u}_{(0)}(\vec{x}, \vec{y}) + \epsilon \vec{u}_{(1)}(\vec{x}, \vec{y}) + \dots$. Finally, the spatial derivatives $\frac{\partial}{\partial x_j}$ within local

434 dynamic equations are replaced by "double scale spatial derivatives" $\frac{\partial}{\partial x_j} + \frac{1}{\epsilon} \frac{\partial}{\partial y_j}$. The factor
 435 in front of the second derivative accounts for the enhancement of the fluctuations at the
 436 local scale.

437 The asymptotic expansion of the displacement field can be introduced into the
 438 dynamic equations (within matrix and inclusion) and continuity equations of displacement
 439 and traction over Γ . Using the double scale derivatives, it produces a set of cell problems
 440 that can be solved sequentially to obtain the displacement fields over matrix and inclusion.

441 The displacement field within the matrix is found to be^{1;2} of the form
 442 $\vec{u}^{(m)}(\vec{x}, \vec{y}) = \vec{U}^{(m)}(\vec{x}) + \epsilon \vec{u}_{(1)}^{(m)}(\vec{x}, \vec{y})$, where $\vec{u}_{(1)}^{(m)}$ is the solution of the cell problem
 443 corresponding to the static homogenization within the matrix for a macroscopic strain
 444 tensor computed from $\vec{U}^{(m)}(\vec{x})$, with null traction over Γ . Indeed, at this stage, due to the
 445 large wavelength within the matrix, the inertia term does not contribute to the localization
 446 problem within the matrix.

447 The first order components $u_{j(0)}^{(i)}$ of the displacement field $\vec{u}_{(0)}^{(i)}$ within the inclusions
 448 are then obtained by solving the first order cell problem for $\vec{u}_{(0)}^{(i)}$:

$$\frac{\partial}{\partial y_k} (\lambda^{(i)} \varepsilon_{u(0)(y)}^{(i)} \delta_{jk} + \mu^{(i)} \varepsilon_{ik(0)(y)}^{(i)}) + \rho^{(i)} \omega^2 u_{j(0)}^{(i)} = 0$$

$$u_{j(0)}^{(i)}(\vec{x}, \vec{y}) = \vec{U}^{(m)}(\vec{x}), \vec{y} \in \Gamma$$

449 where $\varepsilon_{ik(0)(y)}^{(i)}$ stands for the components of the strain tensor within the inclusion computed
 450 from $\vec{u}_{(0)}^{(i)}$ by using y -derivatives.

451 Equation (5) is then obtained by making the change of variable $\vec{w} = \vec{u}_{(0)}^{(i)} - \vec{U}^{(m)}(\vec{x})$

452 **II. DETAILED VALUES OF EIGENFREQUENCIES**

Frequency rank→	1	2	3	4	5	6	7	8
↓Poisson's ratio								
0	2.9345	5.334	7.519	8.561	11.670	12.106	14.861	16.554
0.1	3.050	5.346	7.944	8.590	11.694	12.812	14.865	17.545
0.2	3.230	5.364	8.408	8.828	11.704	13.930	14.876	18.009
0.3	3.553	5.403	8.511	9.981	11.718	14.851	15.977	18.019
0.4	4.277	5.588	8.548	11.678	13.069	14.877	18.013	20.865
0.5	5.136	8.417	11.620	14.796	17.960	21.117	24.270	27.421

Table 1: Values of non-dimensional eigenfrequencies for transversal motion of cylindrical inclusions.

Frequency rank→	1	2	3	4	5	6	7	8
↓Poisson's ratio								
0	3.607	6.129	8.357	9.365	12.444	13.057	15.642	17.538
0.1	3.709	6.148	8.800	9.425	12.472	13.815	15.650	18.553
0.2	3.870	6.178	9.193	9.813	12.488	15.000	15.678	18.789
0.3	4.157	6.241	9.291	11.101	12.518	15.632	17.223	18.807
0.4	4.798	6.494	9.354	12.461	14.521	15.690	18.797	21.908
0.5	5.763	9.095	12.323	15.515	18.689	21.854	25.013	28.168

Table 2: Values of non-dimensional eigenfrequencies for spherical inclusions.

453 **III. * REFERENCES**

- 454 **1.** J.L. Auriault and G. Bonnet, “Dynamique des composites élastiques
455 périodiques(Dynamics of elastic periodic composites)”, Arch. Mech. **37**(4), 269-284
456 (1985).
- 457 **2.** J.L. Auriault, “Acoustics of heterogeneous media: macroscopic behavior by
458 homogenization”, Current Topics Acoust. Res. **1**, 63-90 (1994).
- 459 **3.** J.L. Auriault and C. Boutin, “Long wavelength inner-resonance cut-off frequencies in
460 elastic composite materials”, Int. J. Solids Struct. **49**(23-24), 3269-3281 (2012).

- 461 4. A. Avila, G. Griso and B. Miara, “Phononic bandgaps in linearized elasticity”,
462 Comptes Rendus Acad. Sci. I. **340**(12), 933-938 (2005).
- 463 5. N.O. Babych, I.V. Kamotski and V.P. Smyshlyaev, “Homogenization of spectral
464 problems in bounded domains with doubly high contrasts”, *Networks and*
465 *Heterogeneous Media*. **3**(3), 413-436 (2008).
- 466 6. N.S. Bakhvalov and G.P. Panasenko, *Homogenization averaging process in periodic*
467 *media. Mathematical problems in the mechanics of composite materials* (Kluwer,
468 Dordrecht, 1989), pp.1-366.
- 469 7. A. Baz, “An active acoustic metamaterial with tunable effective density”, *ASME J.*
470 *of Vibration and Acoustics*. **132**(4), 0410111-0410119 (2010).
- 471 8. A. Bensoussan, J.L. Lions and G. Papanicolaou, *Asymptotic analysis for periodic*
472 *structures* (North Holland, Amsterdam, 1978), pp.1-392.
- 473 9. G. Bonnet and V. Monchiet, ” Low frequency locally resonant metamaterials
474 containing composite inclusions”, *J. Acoust. Soc. Am.* **137**(6), 3263-3271 (2015).
- 475 10. J. Carbonell, F. Cervera, J. Sanchez-Dehesa, J. Ariaga, L. Gumen and A. Krokhin,
476 “Homogenization of two-dimensional anisotropic dissipative photonic crystal”, *Appl.*
477 *Phys. Lett.* **97**(4), 231122 (2010).

- 478 **11.** N. Cselyuszka, M. Secujski, V.Crnojevic-Bengin, “Novel negative mass density
479 resonant metamaterial unit cell”, *Phys. Lett. A*, **379**, 33-36 (2015).
- 480 **12.** T.T. Dang, *Méthodes numériques pour l’homogénéisation élastodynamique des*
481 *matériaux hétérogènes périodiques (Numerical Methods for obtaining the dynamic*
482 *effective properties of periodical heterogeneous materials)*, Thesis Université Paris-Est
483 (2015).
- 484 **13.** P. Debye, “Theorie der Spezifischen Wärmen”, *Ann. der Phys.* **39**, 789-839 (1912).
- 485 **14.** Y. Ding, Z. Liu, C. Qiu, and J. Shi, “Metamaterial with Simultaneously Negative
486 Bulk Modulus and Mass Density”, *Phys. Rev. Lett.* **99**, 093904 (2010).
- 487 **15.** A.C. Eringen and E.S. Suhubi, *Elastodynamics, Vol.2* (Acad. press, New York, 1975),
488 pp.772-776 and pp.806-809.
- 489 **16.** R. Gracia-Salgado, D.Torrent and J. Sanchez-Dehesa, “Double-negative acoustic
490 metamaterials based on quasi-two-dimensional fluid-like shells”, *New J. of Phys.*
491 **14**,103052 (2012)
- 492 **17.** H.H. Huang and C.T. Sun, “Locally resonant acoustic metamaterials with 2D
493 anisotropic effective mass density”, *Phil. Mag.* **91**(6),981-996 (2011).
- 494 **18.** H.H. Huang and C.T. Sun, “Anomalous wave propagation in a one-dimensional

- 495 acoustic metamaterial having simultaneously negative mass density and Young's
496 modulus", J. Acous. Soc. Am. **132**(4),2887-2895 (2012).
- 497 **19.** S.H. Lee, C.M. Park, Y.M. Seo, Z.G. Wang and C.K. Kim, "Acoustic metamaterial
498 with negative density", Phys. Lett. A. **373**(48),4464-4469 (2009).
- 499 **20.** Z. Liu, X. Zhang, Y. Mao, Y.Y. Zhu, Z. Yang, C.T. Chan and P. Sheng, "Locally
500 resonant materials", Science **289**,1734-1736 (2000).
- 501 **21.** Z. Liu, C. T. Chan and P. Sheng, "Analytic model of phononic crystals with local
502 resonances", Phys. Rev. B **71**,014103 (2005).
- 503 **22.** A.E.H. Love, *A treatise on the mathematical theory of elasticity* (Dover
504 Publ.,N.Y.,1944), pp. 249-251.
- 505 **23.** G. Ma and P.Sheng, "Acoustic metamaterials: from local resonances to broad
506 horizons", Sci.Adv. **2**,e1501595 (2007).
- 507 **24.** J. Mei, Z. Liu, W. Wen and P. Sheng, "Effective mass density of composites", Phys.
508 Rev. B **76**,134205 (2007).
- 509 **25.** G.W. Milton, "New metamaterials with macroscopic behavior outside that of
510 continuum elastodynamics", New J. Phys. **9**,359 (2007).

- 511 **26.** G.W. Milton and J.R. Willis, “On modifications of Newton’s second law and linear
512 continuum elastodynamics”, Proc. Roy. Soc. A. **463**,855-880 (2007).
- 513 **27.** P.M. Morse and H. Feshbach, *Methods of theoretical physics*. (McGraw-Hill,
514 N.Y.,1953),p. 1264.
- 515 **28.** C.J. Naify, C.M. Chang, G. Mcknight and S.R. Nutt, “Scaling of membrane-type
516 locally resonant acoustic metamaterial arrays”, J. Acoust. Soc. Am.
517 **132**(4),2784-2792 (2012).
- 518 **29.** H Nassar, Q.-C He and N. Auffray. “A generalized theory of elastodynamic
519 homogenization for periodic media”. International Journal of Solids and Structures,
520 **84**,139-146 (2016).
- 521 **30.** H Nassar, Q.-C He and N. Auffray. “On asymptotic elastodynamic homogenization
522 approaches for periodic media”. Journal of the Mechanics and Physics of Solids,
523 **88**,274-290 (2016).
- 524 **31.** S. Nemat-Nasser, J.R. Willis, A. Srivastava and A.V. Amirkhizi, “Homogenization of
525 periodic elastic composites and locally resonant materials”, Phys. Rev. B **83**,104103
526 (2011).

- 527 **32.** A.N. Norris, A.L. Shuvalov and A.A. Kutsenko, “Analytical formulation of
528 three-dimensional dynamic homogenization ”, Proc. R. Soc. A **468**,1629-1651 (2012).
- 529 **33.** J. Park, B. Park, D. Kim and J. Park, “Determination of effective mass density and
530 modulus for resonant metamaterials”, J. Acoust. Soc. Am. **132**(4),2793-2799 (2012).
- 531 **34.** W.J. Parnell and I.D. Abrahams, “Dynamic homogenization in periodic fibre
532 reinforced media. Quasi-static limit for SH waves”, Wave Motion **43**,474-498 (2006).
- 533 **35.** P.J. Schafbuch and F.J. Rizzo, “Eigenfrequencies of an elastic sphere with fixed
534 boundary conditions”, J. Appl. Mech. **59**,458-459 (1992).
- 535 **36.** E. Shamonina and L. Solymar, “Metamaterials: How the subject started”,
536 *Metamaterials*. **1**,12-18 (2007).
- 537 **37.** P. Sheng, *Introduction to wave scattering, localization and mesoscopic phenomena*.
538 (Acad. press, N.Y.,1995),p. 66.
- 539 **38.** P. Sheng, X.X. Zhang, Z. Liu and C.T. Chan, “Locally resonant sonic materials”,
540 *Metamaterials*. **338**,201-205 (2003).
- 541 **39.** J. Soubestre and C. Boutin, “Non-local dynamic behavior of linear fiber reinforced
542 materials”, Mech. Mat. **55**,16-32 (2012).

- 543 **40.** J.O. Vasseur, P.A. Deymier, G. Prantziskonis and G. Hong, “Experimental evidence
544 for the existence of absolute acoustic band gaps in two-dimensional periodic
545 composite media”, *J. Phys.: Condens. Mater.* **10**,6051-6064 (1998).
- 546 **41.** V.G. Veselago, “The electrodynamics of substances with simultaneously negative
547 values of ϵ and μ ”, *Sov. Phys. Usp.* **10**,509-514 (1968).
- 548 **42.** J.R. Willis , “Effective constitutive relation for waves in composites and
549 metamaterials”, *Proc. Roy. Soc. A*, **467**,1865-1879 (2011).
- 550 **43.** Y. Wu, Y. Lai, and Z.-Q. Zhang, “Effective medium theory for elastic metamaterials
551 in two dimensions”, *Phys. Rev. B* **76**,205313 (2007).
- 552 **44.** Z. Yang, J. Mei, M. Yang, N. H. Chan and P. Sheng, “Membrane-Type Acoustic
553 Metamaterial with Negative Dynamic Mass”, *Phys. Rev. Lett.* **101**,204301 (2008).
- 554 **45.** M. Yang, G. Ma, Y. Wu, Z. Yang and P. Sheng, “Homogenization scheme for
555 acoustic metamaterials”, *Phys. Rev. B* **89**,064309 (2014).
- 556 **46.** S. Yao, X. Zhou and G. Hu, “Experimental study on negative effective mass in a 1D
557 mass-spring system”, *New J. Phys.* **10**,043020 (2008).
- 558 **47.** B. Yuan, V. F. Humphrey, J. Wena, X. Wena, “On the coupling of resonance and

559 Bragg scattering effects in three-dimensional locally resonant sonic materials”,

560 Ultrasonics **53**,1332-1343 (2013).

561 **48.** X. Zhou and G. Hu, “Analytic model of elastic metamaterials with local resonances”,

562 Phys. Rev. B **79**,195109 (2009).