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A time-domain spectral finite element method for acoustoelasticity: modeling the effect of mechanical loading on guided wave propagation

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Abstract

Ultrasonic testing techniques such as guided wave-based structural health monitoring aim to evaluate the integrity of a material with sensors and actuators that operate *in situ*, *i.e.* while the material is in use. Since ultrasonic wave propagation is sensitive to environmental conditions such as pre-deformation of the structure, the design and performance evaluation of monitoring systems in this context is a complicated task that requires quantitative data and the associated modeling effort. In our work, we propose a set of numerical tools to solve the problem of mechanical wave propagation in materials subjected to pre-deformation. This type of configuration is usually treated in the domain of acoustoelasticity. A relevant modeling approach is to consider two different problems: a quasi-static nonlinear problem for the large displacement field of the structure and a linearized time-domain wave propagation problem. After carefully reviewing the modeling ingredients to represent the configurations of interest, we propose an original combination of numerical tools that leads to a computationally efficient algorithm. More specifically, we use 3D shell elements for the quasi-static nonlinear problem and the time-domain spectral finite element method to numerically solve the wave propagation problem. Our approach can represent any type of material constitutive law, geometry or mechanical solicitation. We present realistic numerical results on 3D cases related to the monitoring of both isotropic and anisotropic materials, illustrating the genericity and efficiency of our method. We also validate our approach by comparing it to experimental data from the literature.

1 Introduction

In many high-end industries, safety regulations include assessments of the condition of structures and materials used for operations. Typical examples of application areas are nuclear energy, petrochemical industry, transportation, or aeronautics. In fact, any industry where the safety of critical components is of paramount importance. Over time, the need to ensure or assess the health of these components has led to the development of advanced Nondestructive Testing (NDT) techniques such as Structural Health Monitoring (SHM). SHM systems aim to evaluate the structural integrity of the material in question using sensors and processing units *in situ*. In

42 other words, the monitoring systems are attached to the structure once and for all, so
43 that they can continuously generate relevant field data. One means of implementing
44 such systems is to rely on ultrasonic Guided Waves (GW) [50, 65]. These waves have
45 the advantage of propagating over longer distances – compared to bulk waves, which
46 are arguably more common in the NDT community – and thus probe a larger volume
47 of material. Nonetheless, GW propagation is potentially more complex. Namely, it
48 involves dispersion effects that vary from one propagating mode to another, and in
49 the particular case of SHM systems with *in situ* control processes, it is known to
50 be affected by Environmental and Operational Conditions (EOCs) [61]. One of the
51 most general types of EOC that we consider in our work is mechanical solicitations
52 sustained by the structure during its use. These mechanical conditions lead to potential
53 deformations and internal stresses that affect GW propagation. This is the so-called
54 acoustoelastic effect [39, 40]. To provide meaningful assistance in the development
55 and evaluation of SHM system performance, our work addresses the challenge of
56 incorporating acoustoelastic effects into a wave propagation modeling tool that is both
57 generic – *i.e.* applicable to any geometry, constitutive law or mechanical solicitation –
58 and efficient.

59 There are numerous previous works in the literature that address this issue but, to
60 our knowledge, they are mainly based on semi-analytical approaches. For instance, in
61 [41, 48, 51, 52, 53, 56, 55], analytical formulations are proposed and can be solved at
62 low cost. Nevertheless, due to assumptions regarding the geometry or the deformation,
63 such formulations can be limited when precisely modeling complex and realistic cases.
64 For complex shapes and loading configurations, in [32, 42, 57, 60, 62] the authors
65 propose semi-analytical methods that encompass stress inhomogeneity in one or two
66 dimensions. These methods treat the problem in the frequency domain and efficiently
67 compute dispersion curves of guided waves in specified wave guides, given the stress
68 profile through the thickness. In our work, we aim at proposing a time-domain
69 approach for this problem, while considering arbitrary geometries, hyperelastic laws,
70 and inhomogeneous stress profiles caused by mechanical loading.

71 Traditionally, due to the different characteristic times and amplitudes of the
72 structural strain, one can identify two (unidirectionally coupled) problems: (1) a
73 quasi-static nonlinear mechanical problem for the structural deformation; (2) a time-
74 dependent wave problem resulting from a linearization procedure around the quasi-
75 static deformation. In this work, we aim to extend the reach of acoustoelastic GW
76 propagation modeling by combining a set of generic numerical tools to solve both
77 problems. For the first problem, the main challenge (apart from mastering the inherent
78 nonlinearities) is that classical finite element methods (FEMs) give inadequate results
79 for thin, elongated geometries due to numerical locking [30, 38]. To overcome this
80 difficulty we use 3D shell elements [38]. They are based on an asymptotic expansion
81 of the displacement field in the thickness direction and, combined with the so-called
82 Mixed Interpolation of Tensorial Components (MITC) method [14, 20, 24] they are
83 locking-free. Moreover, they are readily available for any kind of constitutive law. The
84 output of this first computational stage is the (large and quasi-static) displacement
85 field sustained by the structure. This output is then used as an input for the wave
86 propagation problem. For this second problem, we use a different type of discretization,
87 namely the Spectral Finite Element Method (SFEM) in the time domain [23, 25, 33].
88 Since the unknown of each problem are fundamentally different in nature, it is natural
89 to resort to different discretization schemes, in particular different meshes. This entails
90 the use of an interpolation operator from the first problem’s mesh onto the second
91 one. The SFEM is based on high-order Lagrangian elements defined on Gauss-Lobatto
92 points. This avoids the Runge phenomenon and achieves spectral-like convergence.

93 Using a consistent mass lumping procedure, they lead to a very efficient explicit
 94 fully discrete scheme – with a stability condition that depends on the mechanical
 95 deformation of the structure. To illustrate the validity and efficiency of our approach,
 96 we provide meaningful 3D simulation results associated with SHM configurations, as
 97 well as experimental validation results with data from the literature [41].

98 This article is divided into two main sections. In *section 2* we explain the
 99 methodology we use in our work. In particular, we first recall the general formalism
 100 of nonlinear elastodynamics, followed by the linearization procedure around a quasi-
 101 static displacement field. In this way, we can accurately define the wave propagation
 102 problem involving the acoustoelastic effects. After comparing our modeling approach
 103 with the relevant previous work on acoustoelasticity, we conclude this section with
 104 the dedicated time and space numerical schemes applied in order to solve the one-
 105 way coupled problems, namely the quasi-static nonlinear problem and the linearized
 106 wave propagation problem. The *section 3* is devoted to the detailed 3D numerical
 107 illustrations of our approach. First, we provide illuminating examples addressing the
 108 potential issue of stability of the fully discrete wave propagation model as a function
 109 of the type of deformation and the constitutive law. Then we provide illustrations
 110 of SHM configurations, namely a tube of isotropic material subjected to a four-point
 111 bending test and a stratified anisotropic plate subjected to torsion. These examples
 112 illustrate the performances and versatility of our approach, while in the last part of
 113 the section we present experimental validations showing the reliability of our modeling
 114 toolchain.

115 2 Methodology for modeling the effect of mechanical load- 116 ing on elastic wave propagation

117 2.1 Nonlinear elastodynamics with quasi-static surface traction & 118 time dependent ultrasonic actuator

119 This section is dedicated to deriving a model representing the propagation of “*low*
 120 *amplitude – high frequency*” waves, within a material subject to a “*high amplitude –*
 121 *(very) low frequency*” mechanical loading. This model enters the traditional framework
 122 of finite deformations of hyperelastic solids. Readers may refer to the reference
 123 textbooks [12, 17, 21] and references therein for an exhaustive presentation of this
 124 generic formalism. We do not seek here to deviate from the very well-understood
 125 formalism of hyperelasticity. Rather, we aim at specifying the corresponding modeling
 126 elements, in particular the various forces that act on the material of interest. This
 127 aspect is of paramount importance through our work, since it leads to a linearization
 128 procedure and its subsequent wave propagation model.

129 Let us denote by $t \geq 0$ the time variable with $T > 0$ the maximal time of interest,
 130 and by $\Omega(t) \subset \mathbb{R}^3$ the volume occupied by the material during its deformation, with
 131 boundary $\Gamma(t)$. We assume that there exists $\Gamma_D \subset \Gamma(t)$ a clamped part of the
 132 boundary. In the following, we consider the functional space $\mathcal{V}(\Omega(t))$ of admissible
 133 virtual displacements vanishing on Γ_D .

134 Quasi-static surface traction.

We consider a sub-part of the boundary $\Gamma_0(t) \subset \Gamma$ that is the area of action of a high
 amplitude and quasi-static force, and such that $\Gamma_0(t) \cap \Gamma_D = \emptyset$. More precisely, let
 us denote by f_0 a prescribed scalar-valued function representing the amplitude of a
 surface traction. The associated virtual work amounts to

$$\int_{\Gamma_0(t)} f_0(\mathbf{x}) \mathbf{n}_0(\mathbf{x}, t) \cdot \mathbf{w}(\mathbf{x}) \, d\Gamma, \quad \forall \mathbf{w} \in \mathcal{V}(\Omega(t)),$$

135 where \mathbf{n}_0 is the outgoing unit normal vector field of Γ_0 . The quasi-static hypothesis
 136 represents the fact that the time variation of this surface traction is assumed to be
 137 significantly lower than the ultrasonic source term described subsequently.

138 Time dependent ultrasonic actuator.

139 We also consider a third part of the boundary $\Gamma_a(t) \subset \Gamma(t)$ where a low amplitude
 140 and time-dependent ultrasonic actuator is vibrating, thus generating propagating wave
 141 packets. We assume that $\Gamma_a(t) \cap \Gamma_0(t) = \emptyset$ and $\Gamma_a(t) \cap \Gamma_D = \emptyset$. We denote by \mathbf{f}_a^δ a
 142 given source term representing this ultrasonic actuator. The superscript δ represents
 143 the amplitude ratio of the ultrasonic source term compared to f_0 . In the following, we
 144 extensively use the assumption that this ratio is small, *i.e.*

$$\delta = \frac{\sup_{(\mathbf{x},t) \in \Omega \times [0;T]} \|\mathbf{f}_a^\delta(\mathbf{x},t)\|_3}{\sup_{\mathbf{x} \in \Omega} |f_0(\mathbf{x})|} \ll 1. \quad (1)$$

In the following, we assume – without loss of generality – that there exists a vector
 field \mathbf{f}_a that corresponds to a rescaled source term, namely $\mathbf{f}_a^\delta = \delta \mathbf{f}_a$. The associated
 virtual work becomes

$$\delta \int_{\Gamma_a(t)} \mathbf{f}_a(\mathbf{x},t) \cdot \mathbf{w}(\mathbf{x}) \, d\Gamma, \quad \forall \mathbf{w} \in \mathcal{V}(\Omega(t)), \quad \forall t \in [0;T].$$

145 Weak formulations of the nonlinear elastodynamic problem.

146 Let us introduce ϱ the mass density of the material. With these notations and the
 147 previous definitions of the various boundary conditions, the principle of virtual works
 148 states that the velocity field $\mathbf{v}(t)$ and the Cauchy stress tensor $\boldsymbol{\sigma}(t)$ satisfy, for any
 149 $\mathbf{w} \in \mathcal{V}(\Omega(t))$

$$\begin{aligned} \int_{\Omega(t)} \varrho(\partial_t \mathbf{v} + (\nabla_x \mathbf{v}) \cdot \mathbf{v}) \cdot \mathbf{w} \, d\Omega + \int_{\Omega(t)} \boldsymbol{\sigma} : \nabla_x \mathbf{w} \, d\Omega \\ = \int_{\Gamma_0(t)} f_0 \mathbf{n}_0 \cdot \mathbf{w} \, d\Gamma + \delta \int_{\Gamma_a(t)} \mathbf{f}_a \cdot \mathbf{w} \, d\Gamma, \end{aligned} \quad (2)$$

150 with given initial conditions $\mathbf{v}|_{t=0}$ and $\boldsymbol{\sigma}|_{t=0}$. Albeit being the standard weak form of
 151 the equilibrium relation, relation (2) deserves some comments. It is not an explicit form
 152 because the deformed domain Ω is unknown, nor is it a closed form since it lacks the
 153 constitutive law linking the Cauchy stress tensor with the kinematic unknown. It should
 154 be noted that this constitutive law is rarely written on a deformed configuration, but
 155 rather on a so-called reference configuration. Also, from a purely practical viewpoint,
 156 the definition of the source term f_a representing the ultrasonic actuator is more likely
 157 to be available on this reference configuration.

From this line of arguments, it is natural to derive a formulation written on a
 reference coordinate system. Let us denote by $\widehat{\Omega}$ the volume occupied by the material
 in this reference setting, and by $\boldsymbol{\xi} \in \widehat{\Omega}$ the associated coordinates. For any time $t > 0$,
 the deformation is represented by the bijective mapping,

$$\begin{aligned} \phi(t) : \widehat{\Omega} &\longmapsto \Omega(t) \\ \boldsymbol{\xi} &\longrightarrow \mathbf{x}(\boldsymbol{\xi},t) = \phi(\boldsymbol{\xi},t), \end{aligned}$$

and we define the displacement field $\mathbf{u}(t)$ as

$$\mathbf{u}(\boldsymbol{\xi},t) = \mathbf{x}(\boldsymbol{\xi},t) - \boldsymbol{\xi}, \quad \forall \boldsymbol{\xi} \in \widehat{\Omega}.$$

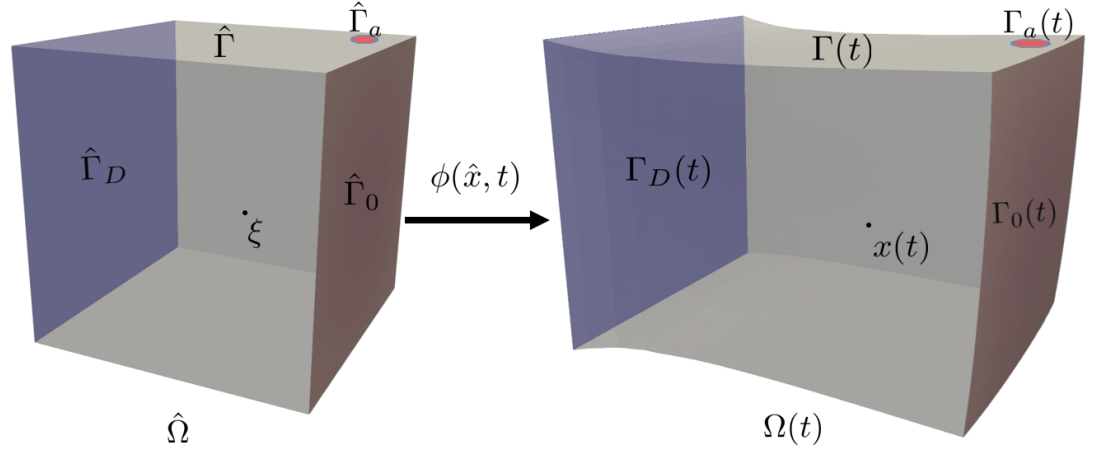


Figure 1. Reference and deformed configurations and their associated notations.

158 We denote by $\hat{\Gamma}$ the boundary of $\hat{\Omega}$, and $\hat{\Gamma}_D$, $\hat{\Gamma}_0$, and $\hat{\Gamma}_a$ the sub-parts of the boundary
 159 corresponding to Γ_D , $\Gamma_0(t)$, and $\Gamma_a(t)$ respectively. See *Figure 1* for an illustration of
 160 both coordinate systems and domains.

The deformation gradient and its determinant are denoted by

$$\mathbf{F}(\boldsymbol{\xi}, t) = \nabla_{\boldsymbol{\xi}}\phi(\boldsymbol{\xi}, t) = \mathbf{I} + \nabla_{\boldsymbol{\xi}}\mathbf{u}(\boldsymbol{\xi}, t) \quad \text{and} \quad J(\boldsymbol{\xi}, t) = \det \mathbf{F}(\boldsymbol{\xi}, t), \quad \forall \boldsymbol{\xi} \in \hat{\Omega},$$

respectively. With these notations, we derive the equivalent of (2) set on the reference configuration, thus explicitly showing the dependency of the displacement field – which is the chosen kinematic unknown in the following. To do so, we apply a standard change of variables from Ω to $\hat{\Omega}$ – see [21] and references therein for more details. Upon the straightforward relation $\mathbf{v} \circ \phi = \partial_t \mathbf{u}$, we start with the inertial terms, namely, for all $\mathbf{w} \in \mathcal{V}(\hat{\Omega})$,

$$\int_{\Omega(t)} \rho(\partial_t \mathbf{v} + (\nabla_{\mathbf{x}} \mathbf{v}) \cdot \mathbf{v}) \cdot (\mathbf{w} \circ \phi^{-1}) \, d\Omega = \int_{\hat{\Omega}} J(\rho \circ \phi) \partial_{tt}^2 \mathbf{u} \cdot \mathbf{w} \, d\hat{\Omega} = \int_{\hat{\Omega}} \hat{\rho} \partial_{tt}^2 \mathbf{u} \cdot \mathbf{w} \, d\hat{\Omega},$$

where $\hat{\rho} = (\rho \circ \phi)J$ is the reference mass density. Concerning the stiffness term, we have, for all $\mathbf{w} \in \mathcal{V}(\hat{\Omega})$,

$$\int_{\Omega(t)} \boldsymbol{\sigma} : \nabla_{\mathbf{x}}(\mathbf{w} \circ \phi^{-1}) \, d\Omega = \int_{\hat{\Omega}} (\boldsymbol{\sigma} \circ \phi) : (\nabla_{\boldsymbol{\xi}} \mathbf{w} \cdot \mathbf{F}^{-1}) J \, d\hat{\Omega} = \int_{\hat{\Omega}} \mathbf{T} : \nabla_{\boldsymbol{\xi}} \mathbf{w} \, d\hat{\Omega},$$

where $\mathbf{T} = (\boldsymbol{\sigma} \circ \phi) \cdot \mathbf{F}^{-\top} J$ is the first Piola-Kirchhoff tensor. Upon the assumption of hyperelasticity [17, 21] there exists a potential function $\mathscr{W}(\mathbf{F})$ that depends exclusively on the deformation gradient and such that

$$\mathbf{T}(\mathbf{F}) = \mathbf{D}_{\mathbf{F}} \mathscr{W}(\mathbf{F}).$$

Note that prescribing a closed form of this energy potential function actually defines the constitutive relation of the material – a matter addressed in the subsequent section. The first Piola-Kirchhoff tensor is non-symmetric, and in the following, we resort to its symmetric counterpart, which can be introduced as follows,

$$\int_{\hat{\Omega}} \mathbf{T} : \nabla_{\boldsymbol{\xi}} \mathbf{w} \, d\hat{\Omega} = \int_{\hat{\Omega}} (\mathbf{F}^{-1} \cdot \mathbf{T}) : (\mathbf{F}^{\top} \cdot \nabla_{\boldsymbol{\xi}} \mathbf{w}) \, d\hat{\Omega} = \int_{\hat{\Omega}} \boldsymbol{\Sigma} : \mathbf{D}_u \mathbf{e}(\mathbf{u}) \mathbf{w} \, d\hat{\Omega},$$

with $\boldsymbol{\Sigma} = \mathbf{F}^{-1} \cdot \mathbf{T}$ the (symmetric) second Piola-Kirchhoff tensor, and

$$\mathbf{e}(\mathbf{u}) = \frac{1}{2} (\nabla_{\boldsymbol{\xi}} \mathbf{u} + \nabla_{\boldsymbol{\xi}} \mathbf{u}^{\top} + \nabla_{\boldsymbol{\xi}} \mathbf{u}^{\top} \cdot \nabla_{\boldsymbol{\xi}} \mathbf{u}), \quad \mathbf{D}_u \mathbf{e}(\mathbf{u}) \mathbf{w} = \frac{1}{2} (\mathbf{F}^{\top} \cdot \nabla_{\boldsymbol{\xi}} \mathbf{w} + \nabla_{\boldsymbol{\xi}} \mathbf{w}^{\top} \cdot \mathbf{F}),$$

the Green-Lagrange tensor and its differential *w.r.t.* the displacement field. For hyperelastic materials, the second Piola-Kirchhoff tensor is usually defined exclusively as a function of the Green-Lagrange tensor, namely $\boldsymbol{\Sigma} = \boldsymbol{\Sigma}(\mathbf{e}(\mathbf{u}))$ in the previous expression, and

$$\boldsymbol{\Sigma}(\mathbf{e}) = \text{D}_e \mathcal{W}(\mathbf{e}).$$

For the quasi-static surface traction, one resorts to the formula for transport of normal vector field to derive the following relations

$$\forall \mathbf{w} \in \mathcal{V}(\widehat{\Omega}), \quad \int_{\Gamma_0(t)} f_0 \mathbf{n}_0 \cdot (\mathbf{w} \circ \phi^{-1}) \, d\Gamma = \int_{\widehat{\Gamma}_0} \widehat{f}_0 (\text{co}\mathbf{F} \cdot \widehat{\mathbf{n}}_0) \cdot \mathbf{w} \, d\widehat{\Gamma},$$

where $\widehat{f}_0 = f_0 \circ \phi$, $\widehat{\mathbf{n}}_0$ is the outgoing unit normal vector field of $\widehat{\Gamma}_0$, and $\text{co}\mathbf{F} = \mathbf{J}\mathbf{F}^{-\top}$ is the co-factor of the deformation gradient. Finally, concerning the ultrasonic source term, we have

$$\forall \mathbf{w} \in \mathcal{V}(\widehat{\Omega}), \quad \int_{\Gamma_a(t)} \mathbf{f}_a \cdot (\mathbf{w} \circ \phi^{-1}) \, d\Gamma = \int_{\widehat{\Gamma}_a} \widehat{\mathbf{f}}_a \cdot \mathbf{w} \, \widehat{\gamma} \, d\widehat{\Gamma},$$

161 where $\widehat{\gamma} = \|\text{co}\mathbf{F} \cdot \widehat{\mathbf{n}}_a\|_3$ with $\widehat{\mathbf{n}}_a$ the unit normal vector field of $\widehat{\Gamma}_a$. This term represents
 162 the effect of the deformation on the surface measure. In the previous relation, $\widehat{\mathbf{f}}_a$ is
 163 the time and space dependency of the actuator on the reference configuration. Its
 164 expression depends on the chosen modeling of the actuator. Without being completely
 165 exhaustive, we can at least provide three main cases of interest:

- 166 • The ultrasonic source is independent of the deformation of the material. For
 167 instance, one could imagine configurations where a hydrophone is emitting
 168 an ultrasonic beam toward the immersed specimen of interest, subject to the
 169 deformation. In that case, we simply have $\widehat{\mathbf{f}}_a = \mathbf{f}_a \circ \phi$.
- The actuator is considered to be generating a time-dependent source term with polarization in the normal direction of the material. This corresponds to the case where there exists a scalar-valued function f_a such that $\mathbf{f}_a = f_a \mathbf{n}_a$, where \mathbf{n}_a is the outgoing unit normal vector field of Γ_a . In that case, we define

$$\widehat{\mathbf{f}}_a = (f_a \circ \phi) \frac{\text{co}\mathbf{F} \cdot \widehat{\mathbf{n}}_a}{\widehat{\gamma}}.$$

- The polarization lies within the local tangent plane of the surface Γ_a . More precisely if we denote by $(\boldsymbol{\tau}_{1,a}, \boldsymbol{\tau}_{2,a})$ two vectors field generating this tangent plane, then there exists two scalar-valued functions $f_{1,a}$ and $f_{2,a}$ such that $\mathbf{f}_a = f_{1,a} \boldsymbol{\tau}_{1,a} + f_{2,a} \boldsymbol{\tau}_{2,a}$, entailing

$$\widehat{\mathbf{f}}_a = \sum_{i=1,2} (f_{i,a} \circ \phi) \frac{\mathbf{F} \cdot \widehat{\boldsymbol{\tau}}_{i,a}}{\|\mathbf{F} \cdot \widehat{\boldsymbol{\tau}}_{i,a}\|_3}, \quad \boldsymbol{\tau}_{i,a} \circ \phi = \mathbf{F} \cdot \widehat{\boldsymbol{\tau}}_{i,a}, \quad \forall i \in \{1, 2\}.$$

170 We can now give the weak formulation of the problem of interest written in the reference
 171 configuration: for any time $t \in [0; T]$, we aim at finding the displacement field $\mathbf{u}(t)$
 172 such that, for any test function $\mathbf{w} \in \mathcal{V}(\widehat{\Omega})$, we have

$$\begin{aligned} \int_{\widehat{\Omega}} \widehat{q} \partial_{tt}^2 \mathbf{u} \cdot \mathbf{w} \, d\widehat{\Omega} + \int_{\widehat{\Omega}} \boldsymbol{\Sigma}(\mathbf{e}(\mathbf{u})) : \text{D}_u \mathbf{e}(\mathbf{u}) \mathbf{w} \, d\widehat{\Omega} \\ = \int_{\widehat{\Gamma}_0} \widehat{f}_0 (\text{co}\mathbf{F} \cdot \widehat{\mathbf{n}}_0) \cdot \mathbf{w} \, d\widehat{\Gamma} + \delta \int_{\widehat{\Gamma}_a} \widehat{\mathbf{f}}_a \cdot \mathbf{w} \, \widehat{\gamma} \, d\widehat{\Gamma}, \end{aligned} \quad (3)$$

173 with given initial conditions $\mathbf{u}|_{t=0}$ and $\partial_t \mathbf{u}|_{t=0}$.

174 2.2 Linearization around a quasi-static displacement field

175 Upon the assumption of a quasi-static mechanical loading and a low amplitude time-
176 dependent ultrasonic source term – as depicted in (1) – we propose to decompose the
177 solution of (3) following the Ansatz

$$\mathbf{u}(\boldsymbol{\xi}, t) = \mathbf{u}_0(\boldsymbol{\xi}) + \delta \mathbf{u}_1(\boldsymbol{\xi}, t) + O(\delta^2), \quad \forall \boldsymbol{\xi} \in \widehat{\Omega}, \quad \forall t \in [0; T]. \quad (4)$$

This (natural) decomposition amounts to seeking a large quasi-static deformation \mathbf{u}_0 generated by the surface traction term, and a low amplitude time-dependent contribution \mathbf{u}_1 emanating from the ultrasonic source term. Intuitively, one would expect \mathbf{u}_0 to satisfy a nonlinear static problem, and \mathbf{u}_1 a wave equation. In this section, we follow standard linearization arguments to derive these two problems. Moreover, following the existence and uniqueness results given in [17], we define the space

$$\mathcal{V}(\widehat{\Omega}) = \left\{ \mathbf{w} \in [H^1(\widehat{\Omega})]^3 \mid \mathbf{w}|_{\Gamma_D} = 0 \right\},$$

178 and the space $\mathcal{W}(\widehat{\Omega})$ of admissible displacements in the reference configuration

$$\mathcal{W}(\widehat{\Omega}) = W^{2,p}(\widehat{\Omega})^3 \cap \mathcal{V}(\widehat{\Omega}), \quad \text{for some } p > 3. \quad (5)$$

179 We now comment briefly this choice. Imbeddings results (see [27, 17]) show that, for
180 any $\mathbf{w}^* \in \mathcal{W}(\widehat{\Omega})$, first, $\nabla_{\boldsymbol{\xi}} \mathbf{w}^*$ belongs to $W^{1,p}(\widehat{\Omega})^{3 \times 3}$ that is an algebra (in particular
181 products of functions in this space do belong to the same space), second, $\nabla_{\boldsymbol{\xi}} \mathbf{w}^*$ is
182 continuous and bounded up to the boundary. We deduce that

$$\mathbf{e}(\mathbf{w}^*) \in W^{1,p}(\widehat{\Omega})^{3 \times 3} \quad \text{and} \quad \forall \mathbf{w} \in \mathcal{V}(\widehat{\Omega}), \quad D_u \mathbf{e}(\mathbf{w}^*) \mathbf{w} \in L^2(\widehat{\Omega})^{3 \times 3}.$$

183 Finally, assuming sufficient smoothness of the potential function $\mathscr{W}(\mathbf{F})$ it can be
184 inferred that $\boldsymbol{\Sigma}(\mathbf{e}(\mathbf{w}^*)) \in W^{1,p}(\widehat{\Omega})^{3 \times 3}$ hence, it is continuous and bounded. A similar
185 smoothness property holds for the fourth order tensor $D_e \boldsymbol{\Sigma}(\mathbf{e}(\mathbf{w}^*))$. The properties
186 given above are implicitly used to give meanings to the integrals introduced in the
187 section below.

188 Stiffness and surface traction operators and associated differentials.

Let us denote by

$$\begin{aligned} \mathbf{A} : \mathcal{W}(\widehat{\Omega}) &\longmapsto \mathcal{V}(\widehat{\Omega})' \\ \mathbf{w}^* &\longrightarrow \mathbf{A}(\mathbf{w}^*), \end{aligned}$$

189 the nonlinear stiffness operator such that, for any $\mathbf{w} \in \mathcal{V}(\widehat{\Omega})$,

$$\langle \mathbf{A}(\mathbf{w}^*), \mathbf{w} \rangle = \int_{\widehat{\Omega}} \boldsymbol{\Sigma}(\mathbf{e}(\mathbf{w}^*)) : D_u \mathbf{e}(\mathbf{w}^*) \mathbf{w} \, d\widehat{\Omega},$$

190 where $\langle \cdot, \cdot \rangle$ denotes the duality product in $\mathcal{V}(\widehat{\Omega})$. The Gâteaux differential of \mathbf{A} at \mathbf{w}^*
191 is denoted by $D_u \mathbf{A}(\mathbf{w}^*) \in \mathcal{L}(\mathcal{V}(\widehat{\Omega}), \mathcal{V}(\widehat{\Omega})')$, and reads, for any $\mathbf{v}, \mathbf{w} \in \mathcal{V}(\widehat{\Omega})$

$$\begin{aligned} \langle D_u \mathbf{A}(\mathbf{w}^*) \mathbf{v}, \mathbf{w} \rangle &= \int_{\widehat{\Omega}} D_u \mathbf{e}(\mathbf{w}^*) \mathbf{v} : D_e \boldsymbol{\Sigma}(\mathbf{e}(\mathbf{w}^*)) : D_u \mathbf{e}(\mathbf{w}^*) \mathbf{w} \, d\widehat{\Omega} \\ &\quad + \int_{\widehat{\Omega}} \boldsymbol{\Sigma}(\mathbf{e}(\mathbf{w}^*)) : D_u^2 \mathbf{e}(\mathbf{v}, \mathbf{w}) \, d\widehat{\Omega}, \quad (6) \end{aligned}$$

where

$$D_u^2 \mathbf{e}(\mathbf{v}, \mathbf{w}) = \frac{1}{2} \left(\nabla_{\boldsymbol{\xi}} \mathbf{v}^{\top} \cdot \nabla_{\boldsymbol{\xi}} \mathbf{w} + \nabla_{\boldsymbol{\xi}} \mathbf{w}^{\top} \cdot \nabla_{\boldsymbol{\xi}} \mathbf{v} \right).$$

192 Note that the differential (6) is obtained from the differential of a product combined
 193 with a chain rule, and the definition of the Green-Lagrange tensor. It actually defines
 194 a bilinear form of particular interest in the following. More precisely, for a given
 195 $\mathbf{w}^* \in \mathcal{W}(\widehat{\Omega})$ defining

$$a(\mathbf{w}^*; \cdot, \cdot) : \mathcal{V}(\widehat{\Omega}) \times \mathcal{V}(\widehat{\Omega}) \mapsto \mathbb{R} \quad (7)$$

$$(\mathbf{v}, \mathbf{w}) \longrightarrow a(\mathbf{w}^*; \mathbf{v}, \mathbf{w}) = \langle \mathbf{D}_u \mathbf{A}(\mathbf{w}^*) \mathbf{v}, \mathbf{w} \rangle,$$

196 then one can verify that $a(\mathbf{w}^*; \cdot, \cdot)$ is a symmetric bilinear form. Note that $\mathbf{D}_u \mathbf{A}$ is
 197 a contribution of two terms: the first one, involving $\mathbf{D}_e \boldsymbol{\Sigma}$, is related to constitutive
 198 nonlinearities since it involves second order derivatives of the potential \mathscr{W} with respect
 199 to the Green-Lagrange tensor \mathbf{e} ; the second one, involving $\boldsymbol{\Sigma}$, is related to geometrical
 200 nonlinearities since it comes exclusively from the linearization of the Green-Lagrange
 201 tensor. Another more compact form – but less prone to interpretation – of the Gâteaux
 202 differential of \mathbf{A} can be deduced from the expression

$$\langle \mathbf{A}(\mathbf{w}^*), \mathbf{w} \rangle = \int_{\widehat{\Omega}} \mathbf{T}(\mathbf{F}(\mathbf{w}^*)) : \nabla_{\xi} \mathbf{w} \, d\widehat{\Omega},$$

203 where $\mathbf{F}(\mathbf{w}^*) = \mathbf{I} + \nabla_{\xi} \mathbf{w}^*$. Computing the Gâteaux differential of \mathbf{A} at \mathbf{w}^* gives the
 204 alternative expression

$$a(\mathbf{w}^*; \mathbf{v}, \mathbf{w}) = \langle \mathbf{D}_u \mathbf{A}(\mathbf{w}^*) \mathbf{v}, \mathbf{w} \rangle = \int_{\widehat{\Omega}} \nabla_{\xi} \mathbf{v} : \mathbf{D}_F \mathbf{T}(\mathbf{F}(\mathbf{w}^*)) : \nabla_{\xi} \mathbf{w} \, d\widehat{\Omega} \quad (8)$$

$$= \int_{\widehat{\Omega}} \nabla_{\xi} \mathbf{v} : \mathbf{D}_F^2 \mathscr{W}(\mathbf{F}(\mathbf{w}^*)) : \nabla_{\xi} \mathbf{w} \, d\widehat{\Omega}.$$

In the following, we also denote by

$$\mathbf{B} : \mathcal{W}(\widehat{\Omega}) \mapsto \mathcal{V}(\widehat{\Omega})'$$

$$\mathbf{w}^* \longrightarrow \mathbf{B}(\mathbf{w}^*),$$

205 the nonlinear surface traction operator such that, for any $\mathbf{w} \in \mathcal{V}(\widehat{\Omega})$,

$$\langle \mathbf{B}(\mathbf{w}^*), \mathbf{w} \rangle = \int_{\widehat{\Gamma}_0} \widehat{f}_0(\text{co}\mathbf{F}(\mathbf{w}^*) \cdot \widehat{\mathbf{n}}_0) \cdot \mathbf{w} \, d\widehat{\Gamma}.$$

206 The Gâteaux differential of \mathbf{B} at \mathbf{w}^* is denoted by $\mathbf{D}_u \mathbf{B}(\mathbf{w}^*) \in \mathcal{L}(\mathcal{V}(\widehat{\Omega}), \mathcal{V}(\widehat{\Omega})')$, and
 207 reads, for any $\mathbf{v}, \mathbf{w} \in \mathcal{V}(\widehat{\Omega})$

$$\langle \mathbf{D}_u \mathbf{B}(\mathbf{w}^*) \mathbf{v}, \mathbf{w} \rangle = \int_{\widehat{\Gamma}_0} \widehat{f}_0((\mathbf{D}_u \text{co}\mathbf{F}(\mathbf{w}^*) \mathbf{v}) \cdot \widehat{\mathbf{n}}_0) \cdot \mathbf{w} \, d\widehat{\Gamma}, \quad (9)$$

where

$$\mathbf{D}_u \text{co}\mathbf{F}(\mathbf{w}^*) \mathbf{v} = \left((\mathbf{F}^{-\top}(\mathbf{w}^*) : \nabla_{\xi} \mathbf{v}) \mathbf{I} - \mathbf{F}^{-\top}(\mathbf{w}^*) \cdot \nabla_{\xi} \mathbf{v} \right) \cdot \text{co}\mathbf{F}(\mathbf{w}^*).$$

Note that the expression of the differential of the cofactor matrix derives from the following relations

$$\mathbf{D}_u \mathbf{J}(\mathbf{w}^*) \mathbf{v} = \text{co}\mathbf{F}(\mathbf{w}^*) : \nabla_{\xi} \mathbf{v}, \quad \mathbf{D}_u (\mathbf{F}^{-\top}(\mathbf{w}^*)) \mathbf{v} = -\mathbf{F}^{-\top}(\mathbf{w}^*) \cdot \nabla_{\xi} \mathbf{v} \cdot \mathbf{F}^{-\top}(\mathbf{w}^*).$$

208 Compared to (6), no particular property – such as symmetry – can be derived from the
 209 bilinear form associated to the differential (9). In order to link with the definition of
 210 the weak formulation (3), we introduce the difference of these two nonlinear operators,
 211 namely

$$\mathbf{A}_b : \mathcal{W}(\widehat{\Omega}) \mapsto \mathcal{V}(\widehat{\Omega})' \quad (10)$$

$$\mathbf{w}^* \longrightarrow \mathbf{A}_b(\mathbf{w}^*) = \mathbf{A}(\mathbf{w}^*) - \mathbf{B}(\mathbf{w}^*),$$

such that its associated Gâteaux differential, for any $\mathbf{w}^* \in \mathcal{W}(\widehat{\Omega})$, simply reads

$$\mathbf{D}_u \mathbf{A}_b(\mathbf{w}^*) = \mathbf{D}_u \mathbf{A}(\mathbf{w}^*) - \mathbf{D}_u \mathbf{B}(\mathbf{w}^*) \in \mathcal{L}(\mathcal{V}(\widehat{\Omega}), \mathcal{V}(\widehat{\Omega})').$$

212 Quasi-static and linearized wave problems.

Let us denote by $\mathcal{H}(\widehat{\Omega}) = [\mathbb{L}^2(\widehat{\Omega})]^3$ and the associated mass-weighted inner product by

$$(\mathbf{v}, \mathbf{w})_{\mathcal{H}} = \int_{\widehat{\Omega}} \widehat{\rho} \mathbf{v} \cdot \mathbf{w} \, d\widehat{\Omega}, \quad \forall \mathbf{v}, \mathbf{w} \in \mathcal{H}.$$

We introduce the nonlinear operator

$$\begin{aligned} L(t, \cdot) : \mathcal{W}(\widehat{\Omega}) &\longmapsto \mathcal{V}(\widehat{\Omega})' \\ \mathbf{w}^* &\longrightarrow L(t, \mathbf{w}^*), \end{aligned}$$

213 such that, for any $\mathbf{w} \in \mathcal{V}(\widehat{\Omega})$,

$$\langle L(t, \mathbf{w}^*), \mathbf{w} \rangle = \int_{\widehat{\Gamma}_a} \widehat{\mathbf{f}}_a(\mathbf{w}^*, t) \cdot \mathbf{w} \widehat{\gamma}(\mathbf{w}^*) \, d\widehat{\Gamma}. \quad (11)$$

From the definition of the operators (10) and (11), we remark that the weak formulation (3) can be written in the following form

$$\frac{d^2}{dt^2}(\mathbf{u}, \mathbf{w})_{\mathcal{H}} + \langle A_b(\mathbf{u}), \mathbf{w} \rangle = \delta \langle L(t, \mathbf{u}), \mathbf{w} \rangle, \quad \forall \mathbf{w} \in \mathcal{V}(\widehat{\Omega}).$$

Using the Ansatz (4), and the time invariance of \mathbf{u}_0 leads to

$$\langle A_b(\mathbf{u}_0), \mathbf{w} \rangle + \delta \left(\frac{d^2}{dt^2}(\mathbf{u}_1, \mathbf{w})_{\mathcal{H}} + \langle D_u A_b(\mathbf{u}_0) \mathbf{u}_1, \mathbf{w} \rangle - \langle L(t, \mathbf{u}_0), \mathbf{w} \rangle \right) = O(\delta^2).$$

214 Hence, satisfying the relation (3) up to second order terms *w.r.t* δ amounts to satisfying
215 the following two problems:

216

217 *i) Find $\mathbf{u}_0 \in \mathcal{W}(\widehat{\Omega})$ the solution, for any $\mathbf{w} \in \mathcal{V}(\widehat{\Omega})$, of*

$$\langle A_b(\mathbf{u}_0), \mathbf{w} \rangle = 0.$$

218 *ii) Find $(0, T) \ni t \mapsto \mathbf{u}_1(t) \in \mathcal{V}(\widehat{\Omega})$ satisfying in $\mathcal{D}'(0, T)$ – a.k.a in the sense of
219 distribution:*

$$\frac{d^2}{dt^2}(\mathbf{u}_1(\cdot), \mathbf{w})_{\mathcal{H}} + \langle D_u A_b(\mathbf{u}_0) \mathbf{u}_1(\cdot), \mathbf{w} \rangle = \langle L(\cdot, \mathbf{u}_0), \mathbf{w} \rangle, \quad \mathbf{w} \in \mathcal{V}(\widehat{\Omega}) \quad (12)$$

220 *with given initial conditions $\mathbf{u}_1(0) \in \mathcal{H}(\widehat{\Omega})$ and $\frac{d\mathbf{u}_1}{dt}(0) \in \mathcal{V}(\widehat{\Omega})'$.*

221 *Approximated linearized wave problem.*

222 A major drawback of the formulation (12) of the wave problem is that it contains
223 the tangent of the surface traction operator, namely $D_u B(\mathbf{u}_0)$, defined in (9). It is
224 associated with a non-symmetric bilinear form, with an undefined sign and implying the
225 surface gradient of its first argument. In addition to these unfortunate mathematical
226 properties, one can interpret this term as the effect of the surface traction onto the wave
227 field. In the applications of interest, where the loading area is restricted compared to
228 the complete surface domain, we expect this term to bear little effect. Combining these
229 arguments, we neglect the contribution of the tangent of the surface traction operator
230 in the wave propagation problem, and we consider the following pair of problems:

231 *i) Find $\mathbf{u}_0 \in \mathcal{V}(\widehat{\Omega})$ the solution, for any $\mathbf{w} \in \mathcal{V}(\widehat{\Omega})$, of*

$$\langle A_b(\mathbf{u}_0), \mathbf{w} \rangle = 0. \quad (13)$$

232 *ii) Find $(0, T) \ni t \mapsto \tilde{\mathbf{u}}(t) \in \mathcal{V}(\widehat{\Omega})$ satisfying in $\mathcal{D}'(0, T)$:*

$$\frac{d^2}{dt^2}(\tilde{\mathbf{u}}(\cdot), \mathbf{w})_{\mathcal{H}} + a(\mathbf{u}_0; \tilde{\mathbf{u}}(\cdot), \mathbf{w}) = \langle L(\cdot, \mathbf{u}_0), \mathbf{w} \rangle, \quad \forall \mathbf{w} \in \mathcal{V}(\widehat{\Omega}) \quad (14)$$

233 *with given initial conditions $\tilde{\mathbf{u}}(0) \in \mathcal{H}((\widehat{\Omega}))$ and $\frac{d\tilde{\mathbf{u}}}{dt}(0) \in \mathcal{V}(\widehat{\Omega})'$, and $a(\mathbf{u}_0; \cdot, \cdot)$*
 234 *defined in (7).*

235 One important remark that can be made at this point is that the two problems
 236 previously defined can be addressed successively. From a practical viewpoint this
 237 means that one can first solve (13), extract the solution \mathbf{u}_0 to compute the tangent
 238 operator and the source term in (14) and solve the wave equation to obtain $\tilde{\mathbf{u}}$.

239 *Well-posedness of the wave propagation problem and energy relation.*

240 Since the focus of our work is on wave propagation modeling, we do not discuss
 241 the questions of the existence and uniqueness of the nonlinear static problem (13).
 242 Theoretical elements on this matter can be found for instance in [12, 17], they usually
 243 entail a so-called “*poly-convexity*” property of \mathcal{W} . We assume in the following that
 244 there exists at least one solution of (13), and rather discuss the well-posedness of the
 245 wave propagation problem of interest (14). We rely on the standard abstract results
 246 that can be found *e.g.* in the reference textbooks [9, 11, 19]. In particular, since the
 247 symmetry of the bilinear form is easily verified from its definition given in (7), we only
 248 need the following assumptions to apply Theorem 4.1 in [11]:

- 249 • *The right-hand side of (14) satisfies $L(\cdot, \mathbf{u}_0) \in \mathcal{C}^1([0, T]; \mathcal{V}(\widehat{\Omega})')$,*
- 250
- 251 • *The bilinear form $a(\mathbf{u}_0; \cdot, \cdot)$ defined in (7) is continuous and coercive, i.e. there*
 252 *exists two strictly positive constants C_0 and α_0 such that*

$$|a(\mathbf{u}_0; \mathbf{v}, \mathbf{w})| \leq C_0 \|\mathbf{v}\|_{\mathcal{V}} \|\mathbf{w}\|_{\mathcal{V}}, \quad a(\mathbf{u}_0; \mathbf{v}, \mathbf{v}) \geq \alpha_0 \|\mathbf{v}\|_{\mathcal{V}}^2, \quad \forall \mathbf{v}, \mathbf{w} \in \mathcal{V}(\widehat{\Omega}). \quad (15)$$

253 Upon these assumptions the problem (14) admits a unique solution such that

$$\tilde{\mathbf{u}} \in L^2((0, T); \mathcal{V}(\widehat{\Omega})), \quad \frac{d\tilde{\mathbf{u}}}{dt} \in L^2((0, T), \mathcal{H}), \quad \frac{d^2\tilde{\mathbf{u}}}{dt^2} \in L^2((0, T), \mathcal{V}(\widehat{\Omega})')$$

254 with additionally, from continuous injection results,

$$\tilde{\mathbf{u}} \in \mathcal{C}^0([0, T]; \mathcal{H}(\widehat{\Omega})) \cap L^\infty([0, T]; \mathcal{V}(\widehat{\Omega})), \quad \frac{d\tilde{\mathbf{u}}}{dt} \in \mathcal{C}^0([0, T]; \mathcal{V}(\widehat{\Omega})') \cap L^\infty([0, T]; \mathcal{H}(\widehat{\Omega})).$$

Note that the coercivity assumption could be replaced by a weaker hypothesis, assuming the coercivity up to a term proportional to $\|\mathbf{v}\|_{\mathcal{H}}^2$. It is clear from the expression of the bilinear form that the coercivity constant depends intricately on the choice of the potential function \mathcal{W} and on the linearization point \mathbf{u}_0 . In particular, in *section 3* we provide illuminating examples of numerical solutions where, depending on the solution of the static problem, we observe a loss of coercivity leading to an ill-posed wave propagation problem. In addition to the previous well-posedness result, we recall the standard energy relation which is an important framework to devise a stable time discretization procedure for the problem (14). Let us denote the kinematic, potential and total energies associated with the unique solution of (14) by

$$\mathcal{K}(t) = \frac{1}{2} \left\| \frac{d\tilde{\mathbf{u}}}{dt} \right\|_{\mathcal{H}}^2, \quad \mathcal{P}(t) = \frac{1}{2} a(\mathbf{u}_0; \tilde{\mathbf{u}}, \tilde{\mathbf{u}}), \quad \mathcal{E} = \mathcal{K} + \mathcal{P}.$$

255 Assuming a null source term in (14), namely $L(t; \mathbf{u}_0) = 0$, then, by using formally
 256 the time derivative of the solution as a test function in (14), we recover the following
 257 energy conservation

$$\frac{d}{dt} \mathcal{E}(t) = 0. \quad (16)$$

258 Energy potentials.

259 Our numerical investigations mainly concern frame-indifferent hyperelastic material
 260 with no residual stress. The material being frame-indifferent the hyperelastic potential
 261 – assumed smooth – is only dependent on the Green-Lagrange tensor \mathbf{e}

$$\mathscr{W}(\mathbf{F}) \equiv \mathscr{W}(\mathbf{e}).$$

262 First, we consider isotropic materials. Since the material has no residual stress, we
 263 necessarily have (see [17]),

$$\mathscr{W}(\mathbf{e}) = \frac{\lambda}{2}(\text{tr } \mathbf{e})^2 + \mu \text{tr } \mathbf{e}^2 + \alpha_1(\text{tr } \mathbf{e})^3 + \alpha_2(\text{tr } \mathbf{e})\text{tr } \mathbf{e}^2 + \alpha_3 \text{tr } \mathbf{e}^3 + o(|\mathbf{e}|^3). \quad (17)$$

264 Note that in the expression above, if the material is heterogeneous, the parameter λ, μ
 265 and the $\{\alpha_i\}_{i=1}^3$ may depend on $\boldsymbol{\xi}$. The coefficients associated with third-order terms,
 266 namely $\{\alpha_i\}_{i=1}^3$, are usually referred to as the Third Order Elastic Constants (TOECs).
 267 The two first terms of this expansion correspond to the so-called Saint-Venant-Kirchhoff
 268 (SVK) potential

$$\mathscr{W}^{\text{SVK}}(\mathbf{e}) = \frac{\lambda}{2}(\text{tr } \mathbf{e})^2 + \mu \text{tr } \mathbf{e}^2,$$

269 while, neglecting higher-order terms, yields the potential

$$\mathscr{W}^{\text{3rd}}(\mathbf{e}) = \mathscr{W}^{\text{SVK}}(\mathbf{e}) + \alpha_1(\text{tr } \mathbf{e})^3 + \alpha_2(\text{tr } \mathbf{e})\text{tr } \mathbf{e}^2 + \alpha_3 \text{tr } \mathbf{e}^3, \quad (18)$$

such potential can be written using Landau's coefficients [15]

$$A = 3\alpha_3, \quad B = \alpha_2 \quad \text{and} \quad C = 3\alpha_1.$$

270 Note that these two potentials are not poly-convex in general, hence stability of the
 271 nonlinear mechanical problem – and therefore of its linearized version – may be an
 272 issue. The most simple elastic potential mitigating this problem is the Compressible
 273 Neo-Hookean (CNH) potential,

$$\mathscr{W}^{\text{CNH}}(\mathbf{e}) = \frac{\lambda}{2}(\sqrt{\det(\mathbf{I} + 2\mathbf{e})} - 1)^2 + \mu \text{tr } \mathbf{e} - \frac{\mu}{2} \log(\det(\mathbf{I} + 2\mathbf{e})).$$

274 The potential \mathscr{W}^{CNH} satisfies the following expansion,

$$\mathscr{W}^{\text{CNH}}(\mathbf{e}) = \frac{\lambda}{2}(\text{tr } \mathbf{e})^2 + \mu \text{tr } \mathbf{e}^2 + \frac{\lambda}{2}(\text{tr } \mathbf{e})^3 - \lambda(\text{tr } \mathbf{e})\text{tr } \mathbf{e}^2 - \frac{4\mu}{3} \text{tr } \mathbf{e}^3 + o(|\mathbf{e}|^3),$$

275 showing that, up to third order term in $|\mathbf{e}|^3$ the CNH constitutive law corresponds to
 276 (18) with $\alpha_1 = \lambda/2$, $\alpha_2 = -\lambda$ and $\alpha_3 = -4\mu/3$. The hyperelastic potential of frame-
 277 indifferent isotropic material can be re-written in terms of the first three invariants of
 278 the Cauchy-Green deformation tensor $\mathbf{C} = \mathbf{F}^T \mathbf{F}$,

$$I_1 = \text{tr}(\mathbf{C}), \quad I_2 = \frac{1}{2}(\text{tr}(\mathbf{C})^2 - \text{tr}(\mathbf{C}^2)), \quad I_3 = \det \mathbf{C}.$$

279 It can be verified that

$$\mathscr{W}^{\text{SVK}} = \frac{(\lambda + 2\mu)}{8} I_1^2 - \frac{(3\lambda + 2\mu)}{4} I_1 - \frac{\mu}{2} I_2 + \frac{(9\lambda + 6\mu)}{8} \quad (19)$$

280 and

$$\mathscr{W}^{\text{CNH}} = \frac{\lambda}{2}(\sqrt{I_3} - 1)^2 + \frac{\mu}{2}(I_1 - 3 - \log(I_3)). \quad (20)$$

281 Murnaghan [2] proposed different TOECs to characterize the second-order elastic
282 deformations, namely l , m and n ,

$$\mathscr{W}^{3\text{rd}} = \mathscr{W}^{\text{SVK}} + \frac{l}{24}(I_1 - 3)^3 + \frac{m}{12}(I_1 - 3)(I_1^2 - 3I_2) + \frac{n}{8}(I_1 - I_2 + I_3 - 1). \quad (21)$$

283 In the context of acoustoelasticity, Hughes and Kelly related these TOECs with changes
284 in bulk wave velocities for prestressed isotropic materials and experimentally assessed
285 them in [3]. The relations between (l, m, n) and $(\alpha_1, \alpha_2, \alpha_3)$ appearing in the expression
286 (18) are given by

$$\alpha_1 = l - \frac{m}{3} + \frac{n}{6}, \quad \alpha_2 = m - \frac{n}{2}, \quad \alpha_3 = \frac{n}{3}.$$

287 We also investigate an orthotropic configuration where one plane of symmetry is
288 lost, orthogonal to a principal axis a , in that case, the material is said to be transversely
289 isotropic. To fully describe its constitutive behavior we need to add two more invariants

$$I_4 = a \cdot \mathbf{C} \cdot a, \quad I_5 = a \cdot \mathbf{C}^2 \cdot a.$$

290 One example of the hyperelastic potential for transversely isotropic materials is the
291 one from [22],

$$\mathscr{W}^{\text{CNH-TI}} = \mathscr{W}^{\text{CNH}} + [\alpha + \beta \ln J + \gamma(I_4 - 1)](I_4 - 1) - \frac{\alpha}{2}(I_5 - 1), \quad (22)$$

292 where α , β and γ are additional parameters that can be calibrated to match observed
293 transversely isotropic behavior.

294 *Acoustoelastic effects.*

295 The change in the behavior of elastic waves when propagating in a deformed or
296 prestressed medium is known as acoustoelastic effect. This phenomenon is extensively
297 studied, see *e.g.* [1, 2, 3]. In [13], the author analyses the structure and properties of
298 incremental elastic deformations models, including the unicity and stability aspects
299 with respect to the constitutive law. As far as isotropic materials are considered, the
300 behavior of elastic waves – bulk wave velocities, in particular – has been mathematically
301 studied by the authors in [39]. We can relate their considerations to stability properties
302 and positivity of the bilinear form $a(\mathbf{u}_0; \cdot, \cdot)$. Indeed, from the expression (8), assuming
303 that the solution behaves locally as a plane wave, with polarization \mathbf{m} propagating in
304 the direction \mathbf{n} , and assuming smoothness of the parameters and of \mathbf{u}_0 , we arrive at
305 the necessary condition for the positivity of $a(\mathbf{u}_0; \cdot, \cdot)$,

$$\forall (\mathbf{m}, \mathbf{n}) \in \mathbb{R}^3 \times \mathbb{R}^3, \quad \mathbf{m} \otimes \mathbf{n} : \mathbb{D}_{\mathbf{F}}^2 \mathscr{W}(\mathbf{F}(\mathbf{u}_0)) : \mathbf{m} \otimes \mathbf{n} \geq 0. \quad (23)$$

306 Note that the inequality above should be verified for each $\boldsymbol{\xi} \in \widehat{\Omega}$ since \mathbf{u}_0 is a function
307 of $\boldsymbol{\xi}$ even when the material is homogeneous. Now, following [39], this amounts to
308 checking the positivity of the so-called (symmetric) Christoffel tensor $\mathbf{Q}(\mathbf{u}_0, \mathbf{n}) \in \mathbb{R}^{3 \times 3}$,
309 indeed,

$$\forall (\mathbf{m}, \mathbf{n}) \in \mathbb{R}^3 \times \mathbb{R}^3, \quad \mathbf{m} \otimes \mathbf{n} : \mathbb{D}_{\mathbf{F}}^2 \mathscr{W}(\mathbf{F}(\mathbf{u}_0)) : \mathbf{m} \otimes \mathbf{n} = \mathbf{m} \cdot \mathbf{Q}(\mathbf{u}_0, \mathbf{n}) \cdot \mathbf{m}.$$

310 To further analyse it, we use the Christoffel tensor $\mathbf{Q}(\mathbf{u}_0, \mathbf{n})$ given in [39] in the
311 isotropic case and for the reference configuration. It depends on \mathscr{W}_i and \mathscr{W}_{ij} , the first
312 and second-order derivatives of the potential with respect to the invariant I_i and I_j
313 respectively. Note that the dependence in \mathbf{u}_0 of \mathscr{W}_i and \mathscr{W}_{ij} is temporarily omitted for
314 the sake of clarity. In the case where

$$\mathscr{W}_{12} = \mathscr{W}_{22} = \mathscr{W}_{13} = \mathscr{W}_{23} = 0,$$

315 which is verified for the Saint-Venant-Kirchhoff and Compressible Neo-Hookean consti-
316 tutive laws, the Christoffel tensor is given by

$$\begin{aligned} \mathbf{Q}(\mathbf{u}_0, \mathbf{n}) &= (4I_3^2 \mathcal{W}_{33} + 2I_3 \mathcal{W}_3) \hat{\mathbf{n}} \otimes \hat{\mathbf{n}} + 2(\mathcal{W}_1 + I_1 \mathcal{W}_2) \mathbf{I} \\ &+ (4\mathcal{W}_{11} + 2\mathcal{W}_2) (\mathbf{F}(\mathbf{u}_0) \cdot \mathbf{n}) \otimes (\mathbf{F}(\mathbf{u}_0) \cdot \mathbf{n}) - 2\mathcal{W}_2 \left(\|\mathbf{F}(\mathbf{u}_0) \cdot \mathbf{n}\|_3^2 \mathbf{I} + \mathbf{B}(\mathbf{u}_0) \right), \end{aligned} \quad (24)$$

317 where $\mathbf{B}(\mathbf{u}_0) = \mathbf{F}(\mathbf{u}_0) \mathbf{F}^\top(\mathbf{u}_0)$ and $\hat{\mathbf{n}} = \mathbf{F}^{-\top}(\mathbf{u}_0) \cdot \mathbf{n}$. The expression above takes a
318 simple form when considering the CNH constitutive law

$$\mathbf{Q}^{\text{CNH}}(\mathbf{u}_0, \mathbf{n}) = (\lambda I_3(\mathbf{u}_0) + \mu) \hat{\mathbf{n}} \otimes \hat{\mathbf{n}} + \mu \mathbf{I}.$$

319 When considering the SVK constitutive law, we get

$$\begin{aligned} \mathbf{Q}^{\text{SVK}}(\mathbf{u}_0, \mathbf{n}) &= \frac{1}{2} (I_1(\mathbf{u}_0) \lambda - 3\lambda - 2\mu + 2\mu \|\mathbf{F}(\mathbf{u}_0) \cdot \mathbf{n}\|^2) \mathbf{I} \\ &+ (\lambda + \mu) (\mathbf{F}(\mathbf{u}_0) \cdot \mathbf{n}) \otimes (\mathbf{F}(\mathbf{u}_0) \cdot \mathbf{n}) + \mu \mathbf{B}(\mathbf{u}_0). \end{aligned} \quad (25)$$

320 Considering pure dilatation deformation, *i.e.*, $\mathbf{u}_0 = \boldsymbol{\xi}(d_0 - 1)$ and $\mathbf{F}(\mathbf{u}_0) = d_0 \mathbf{I}$ with
321 $d_0 \geq 1$ we have,

$$\mathbf{Q}^{\text{SVK}}(\boldsymbol{\xi}(d_0 - 1), \mathbf{n}) = \frac{(d_0^2 - 1)}{2} (3\lambda + 2\mu) \mathbf{I} + d_0^2 [(\lambda + \mu) \mathbf{n} \otimes \mathbf{n} + \mu \mathbf{I}]$$

322 and

$$\mathbf{Q}^{\text{CNH}}(\boldsymbol{\xi}(d_0 - 1), \mathbf{n}) = (\lambda d_0^4 + \mu d_0^{-2}) \mathbf{n} \otimes \mathbf{n} + \mu \mathbf{I},$$

323 One can see that the eigenvalues of $\mathbf{Q}^{\text{SVK}}(\boldsymbol{\xi}(d_0 - 1), \mathbf{n})$, are increasing functions of d_0 ,
324 meaning that waves tend to propagate faster – at least locally. The CNH constitutive
325 law behaves similarly when $2\lambda \geq \mu$. This behavior contradicts observed experiments
326 [3] where velocities of the shear and pressure waves decrease for some materials –
327 aluminum typically – that undergoes an increase in volume. In this case, proper
328 behavior can be modeled using high-order elasticity models such as the Murnaghan's,
329 despite some expected and observed stability issues (see the numerical experiments
330 after). In an hypothetical case of pure stretching in one direction, we have locally

$$\mathbf{F}(\mathbf{u}_0) = \begin{pmatrix} d_0 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & 1 \end{pmatrix},$$

331 that is to say $\mathbf{u}_0 = \mathbf{e}_1(\mathbf{e}_1 \cdot \boldsymbol{\xi})(d_0 - 1)$, with $d_0 \geq 1$, we have

$$\mathbf{Q}^{\text{SVK}}(\mathbf{u}_0, \mathbf{e}_1) = \left(\frac{d_0^2 - 1}{2} (\lambda + 2\mu) + \mu \right) \mathbf{I} + (d_0^2 (\lambda + 2\mu) - \mu) \mathbf{e}_1 \otimes \mathbf{e}_1,$$

332 and

$$\mathbf{Q}^{\text{CNH}}(\mathbf{u}_0, \mathbf{e}_1) = (\lambda + d_0^{-2} \mu) \mathbf{e}_1 \otimes \mathbf{e}_1 + \mu \mathbf{I}.$$

333 In this particular case, the behavior of SVK and CNH constitutive laws differs. Indeed
334 one can observe that for increasing values of d_0 the eigenvalues of \mathbf{Q}^{SVK} are all
335 increasing while the eigenvalues of \mathbf{Q}^{CNH} associated to the eigenvector \mathbf{e}_1 is given by
336 $\lambda + (d_0^{-2} + 1)\mu$, hence it decreases with d_0 . This behavior is observed in our numerical
337 results.

338 In the acoustoelasticity literature, a more general formulation comprising anisotropic
 339 behavior is usually presented. The expansion of the potential up to third-order terms
 340 are written in the indicial notation as

$$\mathcal{W}(e) \approx C_{ijkl}e_{ij}e_{kl} + C_{ijklmn}e_{ij}e_{kl}e_{mn}, \quad (26)$$

341 which is a generalization of the expansion (17). The fourth-order tensor C_{ijkl} is
 342 the traditional elasticity tensor, with components related to the *second order elastic*
 343 *constants* from Hooke's law, and the sixth-order tensor C_{ijklmn} has its components
 344 related to the *third-order elastic constants*. One may note that this description of the
 345 constitutive behavior complexifies significantly the study of poly-convexity.

346 Obtaining material parameters for suitable constitutive models is an essential and
 347 difficult task. Many works concern the experimental assessment of these parameters.
 348 For example, the seminal work presented in [3] gives the TOECs for polystyrene, iron,
 349 and Pyrex glass. In [4, 5, 6, 8, 7, 10, 26, 36, 41, 44, 45], the TOECs are obtained for
 350 aluminum and other metals. Transversely isotropic materials are specifically treated
 351 in [18, 46, 49, 53]. Most of the presented methods rely on pitch-catch time of flight
 352 measurements, which is potentially subject to poor robustness [36]. Alternatively,
 353 more sophisticated methods based on laser vibrometer aided measurements and space-
 354 time Fourier Transform can be applied to enhance the quality of the parameteres
 355 reconstruction [59]. Note that, in practice, it appears that experimental set up
 356 uncertainties can have a large influence on the estimation of the parameters, as studied
 357 in [36]. Additionally, fabrication processes may significantly modify the acoustoelastic
 358 response of the material [10, 16].

359 2.3 Dedicated time & space numerical schemes

360 We now move on to proposing a discretization strategy that addresses the problem
 361 of solving the one-way coupled problems “ $i \Rightarrow ii$ ” defined in (13) and (14). Each
 362 problem, namely the nonlinear static problem and the wave propagation problem,
 363 presents specific numerical challenges – that have been tackled independently in the
 364 literature – and we propose here to aggregate the suitable numerical methods to derive
 365 a complete and efficient fully discrete algorithm.

366 Solving the quasi-static problem using 3D shell elements.

The nonlinearities associated to problem (13) are handled using Newton-Raphson
 iterations: for a given \mathbf{u}_0^k , find $\mathbf{u}_0^{k+1} \in \mathcal{W}(\widehat{\Omega})$ such that

$$\langle \mathbf{D}_u A_b(\mathbf{u}_0^k)(\mathbf{u}_0^{k+1} - \mathbf{u}_0^k), \mathbf{w} \rangle = \langle A_b(\mathbf{u}_0^k), \mathbf{w} \rangle, \quad \forall \mathbf{w} \in \mathcal{V}(\widehat{\Omega}),$$

367 until convergence. In some cases, typically when the maximal amplitude of the surface
 368 traction is large, this iterative process may not converge. To increase robustness
 369 in these cases, we apply a continuation algorithm – *i.e.* successive reductions of
 370 the loading amplitude, see [47] and references therein for more details. Writing the
 371 increment $\mathbf{u} = \mathbf{u}_0^{k+1} - \mathbf{u}_0^k$, we need to solve at each iteration a linearized elasticity
 372 problem in the form of

373 Find $\mathbf{u} \in \mathcal{W}(\widehat{\Omega})$ the solution, for any $\mathbf{w} \in \mathcal{V}(\widehat{\Omega})$ of

$$a_k(\mathbf{u}, \mathbf{w}) - b_k(\mathbf{u}, \mathbf{w}) = \ell_k(\mathbf{w}), \quad (27)$$

where $a_k(\cdot, \cdot) = a(\mathbf{u}_0^k; \cdot, \cdot)$ defined in (7), and

$$b_k(\mathbf{u}, \mathbf{w}) = \langle \mathbf{D}_u \mathbf{B}(\mathbf{u}_0^k)\mathbf{u}, \mathbf{w} \rangle, \quad \ell_k(\mathbf{w}) = \langle A_b(\mathbf{u}_0^k), \mathbf{w} \rangle.$$

374 It is well-known that for configurations where $\widehat{\Omega}$ is a thin elongated geometry, standard
 375 finite element procedures have very poor performances, or can even fail to provide
 376 meaningful results due to locking phenomena [30, 38]. To circumvent this difficulty
 377 we use 3D shell elements [38] that enter the framework of general shell elements [14,
 378 20]. These elements possess desirable properties in the scope of our work: (1) they are
 379 locking-free; (2) they rely on the full 3D variational formulations of the type of (27),
 380 hence, they can be applied to any type of constitutive law; (3) they rely on a quadratic
 381 expansion of the displacement unknown, involving nodes on the outer surface. This
 382 makes them readily suitable for coupling with other elements or with themselves, *i.e.*
 383 stacking of shell elements. This aspect is particularly useful to represent stratified
 materials.

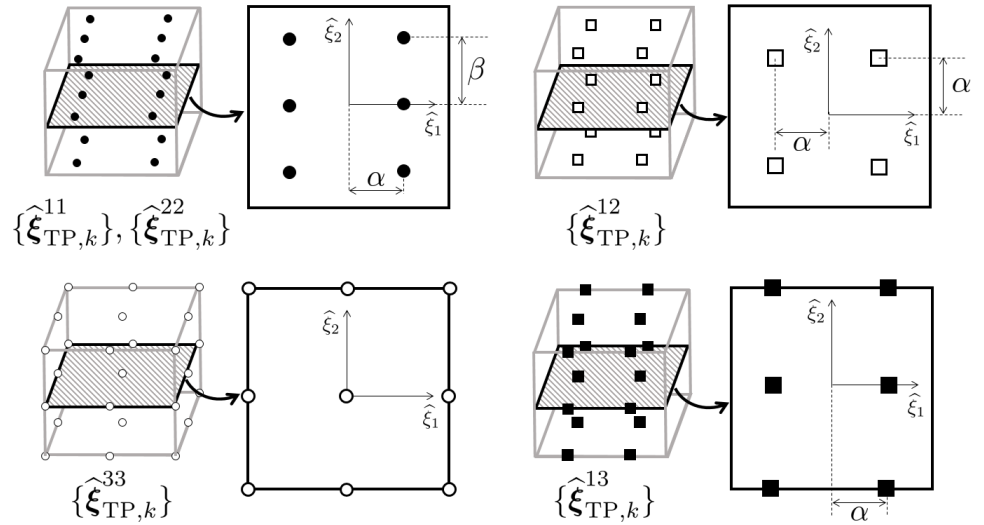


Figure 2. Definition of the MITC9 tying points on the reference hexahedron – extracted from [24] – with $\alpha = 1/\sqrt{3}$ and $\beta = \sqrt{3}/\sqrt{5}$. The tying points for the remaining tensorial components are deduced by symmetry.

384

Let us assume for illustration purposes that the material at hand is homogeneous and that the domain of interest can be expressed as $\widehat{\Omega} = \widehat{\omega} \times]-\eta; \eta[$, where $\widehat{\omega} \subset \mathbb{R}^2$ is the mid-surface, and $\eta > 0$ is the half-thickness of the domain. Since 3D shell elements rely on a quadratic expansion of the displacement unknown, after defining quadrangular quadratic finite elements on the mid-surface, they amount to considering standard Q^2 Lagrange elements on the full 3D domain [38]. More precisely, let us define \mathcal{T}_H , a tessellation of $\widehat{\Omega}$ with mesh step H and composed of hexahedra with one element in the thickness. Denoting by $\widehat{K} = [-1; 1]^3$ the reference element, any hexahedron $K \in \mathcal{T}_H$ is the deformation of the reference element through a \mathcal{C}^1 -diffeomorphism denoted by \mathcal{F}_K . We consider the following approximation space

$$V_H(\widehat{\Omega}) = \left\{ w_H \in \mathcal{C}^0(\widehat{\Omega}) \mid \forall K \in \mathcal{T}_H, \exists \widehat{w} \in Q^2(\widehat{K}), w_H|_K = \widehat{w} \circ \mathcal{F}_K^{-1} \right\},$$

and its vectorial counterpart including the relevant boundary conditions

$$\mathcal{V}_H(\widehat{\Omega}) = \left\{ \mathbf{w}_H \in [V_H(\widehat{\Omega})]^3, \mathbf{w}_H|_{\Gamma_D} = 0 \right\} \subset \mathcal{V}(\widehat{\Omega}).$$

Considering a direct application of an approximation of problem (27) using $\mathcal{V}_H(\widehat{\Omega})$ would not alleviate the locking phenomena. To do so, one can resort to the so-called Mixed Interpolation of Tensorial Components (MITC) method [14, 20, 24]. In essence, this approach consists of interpolating the (discrete) Green-Lagrange tensor on so-called

tying points. The location of these points is different from the Lagrange interpolation points defining the discrete displacement unknown and is also different from one component of the interpolated tensor to another. In our work, we use the MITC9 point distribution since we consider quadratic elements – see *Figure 2* for the positioning of these points. Let $i, j \in \{1, 2, 3\}$ be indexes of the components of the Green-Lagrange tensor, and n_{TP}^{ij} the total number of the tying points denoted by

$$\{\widehat{\boldsymbol{\xi}}_{\text{TP},k}^{ij}\}_{k=1}^{n_{\text{TP}}^{ij}} \subset \widehat{K},$$

and defined on the reference hexahedron. The MITC process applied to a component e_{ij} of the Green-Lagrange tensor is denoted by $\mathcal{I}_{\text{TP}}(e_{ij})$ and reads

$$\forall K \in \mathcal{T}_H, \quad \forall \boldsymbol{\xi} \in K, \quad \mathcal{I}_{\text{TP}}(e_{ij})|_K(\boldsymbol{\xi}) = \sum_{k=1}^{n_{\text{TP}}^{ij}} \widehat{\lambda}_{\text{TP},k}(\mathcal{F}_K^{-1}(\boldsymbol{\xi})) e_{ij}(\mathcal{F}_K(\widehat{\boldsymbol{\xi}}_{\text{TP},k}^{ij})),$$

where $\{\widehat{\lambda}_{\text{TP},i}\}_{i=1}^{n_{\text{TP}}}$ are the Lagrange shape functions associated with the tying points. In the following, we denote by $\bar{\mathbf{e}} = \mathcal{I}_{\text{TP}}(\mathbf{e})$ the MITC process applied to tensors, in the sense that each component is interpolated on the tying points, namely

$$\forall i, j \in \{1, 2, 3\}, \quad \bar{e}_{ij} = \mathcal{I}_{\text{TP}}(e)_{ij} = \mathcal{I}_{\text{TP}}(e_{ij}).$$

In the same fashion, considering a discrete linearization point $\mathbf{w}_H^* \in \mathcal{V}_H(\widehat{\Omega})$, we introduce, for any $\mathbf{v}_H, \mathbf{w}_H \in \mathcal{V}_H(\widehat{\Omega})$,

$$\mathbf{D}_u \bar{\mathbf{e}}(\mathbf{w}_H^*) \cdot \mathbf{w}_H = \mathcal{I}_{\text{TP}}(\mathbf{D}_u \mathbf{e}(\mathbf{w}_H^*) \cdot \mathbf{w}_H), \quad \mathbf{D}_u^2 \bar{\mathbf{e}}(\mathbf{v}_H, \mathbf{w}_H) = \mathcal{I}_{\text{TP}}(\mathbf{D}_u^2 \mathbf{e}(\mathbf{v}_H, \mathbf{w}_H)),$$

385 the interpolation on the tying points of the first and second differential of the Green-
386 Lagrange tensor. With these notations, the discrete problem associated with (27)
387 reads:

388 Find $\mathbf{u}_H \in \mathcal{V}_H(\widehat{\Omega})$ the solution, for any $\mathbf{w}_H \in \mathcal{V}_H(\widehat{\Omega})$, of

$$\bar{a}_k(\mathbf{u}_H, \mathbf{w}_H) - b_k(\mathbf{u}_H, \mathbf{w}_H) = \ell_k(\mathbf{w}_H),$$

389 with

$$\begin{aligned} \bar{a}_k(\mathbf{u}_H, \mathbf{w}_H) &= \int_{\widehat{\Omega}} \mathbf{D}_u \bar{\mathbf{e}}(\mathbf{u}_{0,H}^k) \mathbf{u}_H : \mathbf{D}_e \boldsymbol{\Sigma}(\bar{\mathbf{e}}(\mathbf{u}_{0,H}^k)) : \mathbf{D}_u \bar{\mathbf{e}}(\mathbf{u}_{0,H}^k) \mathbf{w}_H \, d\widehat{\Omega} \\ &\quad + \int_{\widehat{\Omega}} \boldsymbol{\Sigma}(\bar{\mathbf{e}}(\mathbf{u}_{0,H}^k)) : \mathbf{D}_u^2 \bar{\mathbf{e}}(\mathbf{u}_H, \mathbf{w}_H) \, d\widehat{\Omega}. \end{aligned}$$

390 In the previous expression of the discrete problem, $\mathbf{u}_{0,H}^k \in \mathcal{V}_H(\widehat{\Omega})$ represents the
391 discrete displacement field obtained from the prior Newton-Raphson iteration, and is
392 also used as a linearization point for the bilinear and linear forms $b_k(\cdot, \cdot)$ and $\ell_k(\cdot)$.

393 Spectral elements and explicit leap-frog scheme for the wave propagation problem.
394 We now provide key elements for efficient space and time discretization of the wave
395 problem (14) relying on the combination of spectral finite elements, the mass-lumping
396 technique and the explicit leapfrog scheme. Readers may refer to [23, 25, 33] for a
397 more detailed description of these subjects.

To start with, let us denote by $\mathbf{p} = (p_1, p_2, p_3)^\top \in \mathbb{N}^{*3}$ an integer triplet representing orders of polynomial on every direction of the reference hexahedron. We define the space of polynomial with “anisotropic orders” as

$$Q^{\mathbf{p}}(\widehat{K}) = \text{span} \{ \widehat{s}_1^i \widehat{s}_2^j \widehat{s}_3^k, \quad \forall (i, j, k) \in \mathbb{N} : i \leq p_1, j \leq p_2, k \leq p_3 \},$$

where $\widehat{\mathbf{s}} = (\widehat{s}_1, \widehat{s}_2, \widehat{s}_3)^\top$ are the local coordinates on the reference element. From this local space, we can build the global conforming approximation space that supports the semi-discretization procedure of the wave problem (14). To do so, we consider \mathcal{T}_h a hexahedral mesh of $\widehat{\Omega}$ – different from \mathcal{T}_H – with mesh step h and such that any hexahedron $K \in \mathcal{T}_h$ is the deformation of the reference element through a \mathcal{C}^1 -diffeomorphism denoted by \mathcal{G}_K . The scalar-valued approximation space reads

$$V_h(\widehat{\Omega}) = \left\{ w_h \in \mathcal{C}^0(\widehat{\Omega}) \mid \forall K \in \mathcal{T}_h, \exists \widehat{w} \in Q^p(\widehat{K}), w_h|_K = \widehat{w} \circ \mathcal{G}_K^{-1} \right\},$$

and, as previously, its vectorial counterpart including relevant boundary conditions is

$$\mathcal{V}_h(\widehat{\Omega}) = \left\{ \mathbf{w}_h \in [V_h(\widehat{\Omega})]^3, \mathbf{w}_h|_{\Gamma_D} = 0 \right\} \subset \mathcal{V}(\widehat{\Omega}).$$

398 The local Lagrange (polynomial) basis functions are defined on the reference element
 399 from Gauss-Lobatto points. This is a key point in order to avoid Runge’s phenomenon,
 400 and to render spectral-like convergence as the order of the local polynomials increases.
 401 In practice, these points are set on the reference edge $\widehat{e} = [-1; 1]$ for any order of
 402 approximation (see *e.g.* [33] for an explicit definition and related examples of these
 403 points), and the location of points are deduced in the reference cube by a simple
 404 extrusion process.

In addition to this specific choice of approximation space, we use a quadrature formula to numerically compute the integrals involved in the bilinear forms $(\cdot, \cdot)_{\mathcal{H}}$ and $a(\mathbf{u}_0; \cdot, \cdot)$ appearing in (14). The formula is defined on the reference hexahedron. We denote the \widehat{n}_q associated weights and points respectively by

$$(\{\widehat{\omega}_i\}_{i=1}^{\widehat{n}_q}, \{\widehat{\mathbf{s}}_i\}_{i=1}^{\widehat{n}_q}).$$

For any function f at least continuous over each element in \mathcal{T}_h , the numerical integration process reads

$$\int_{\Omega} f \, d\Omega \approx \sum_{K \in \mathcal{T}_h} \sum_{k=1}^{\widehat{n}_q} \widehat{\omega}_k f(\mathcal{G}_K(\widehat{\mathbf{s}}_k)) |\det \nabla_{\widehat{\mathbf{s}}} \mathcal{G}_K(\widehat{\mathbf{s}}_k)|.$$

405 The quadrature points are identical to the Gauss-Lobatto points, and the quadrature
 406 weights are the strictly positive Gauss-Lobatto weights. Matching the quadrature
 407 points with the points associated with the local Lagrange basis functions is referred to
 408 as “mass lumping” [23, 25, 33]. It leads to a diagonal approximated mass matrix. It
 409 has been the subject of particular interest in the context of finite element methods
 410 applied to time domain wave propagation problems since it enables fully explicit time
 411 discretization schemes. To conserve stability and accuracy, specific conditions must be
 412 enforced on the quadrature points and weights. For quadrangles and hexahedra, these
 413 conditions are met at any order of approximations by the Gauss-Lobatto rules [33, 35].
 414 Following this numerical integration procedure, we denote by

$$(\cdot, \cdot)_{\mathcal{H}} \approx (\cdot, \cdot)_h, \quad a(\mathbf{u}_0; \cdot, \cdot) \approx a_h(\mathbf{u}_0; \cdot, \cdot), \quad (28)$$

415 the approximated bilinear forms so that the semi-discrete formulation of the wave
 416 problem (14) reads

417 For any time $t \in [0; T]$, find $\widetilde{\mathbf{u}}_h(t) \in \mathcal{V}_h(\widehat{\Omega})$ the solution, for any $\mathbf{w}_h \in \mathcal{V}_h(\widehat{\Omega})$, of

$$\frac{d^2}{dt^2} (\widetilde{\mathbf{u}}_h, \mathbf{w}_h)_h + a_h(\mathbf{u}_{0,H}; \widetilde{\mathbf{u}}_h, \mathbf{w}_h) = \langle L(t, \mathbf{u}_{0,H}), \mathbf{w}_h \rangle. \quad (29)$$

418 with given initial conditions $\widetilde{\mathbf{u}}_h|_{t=0}$ and $\partial_t \widetilde{\mathbf{u}}_h|_{t=0}$.

419 Note that in (29) $\mathbf{u}_{0,H}$ is the discrete solution – after convergence of the Newton-
 420 Raphson iterations – of problem (13). Moving on to the time discretization of the wave
 421 propagation problem, we define $\Delta t > 0$ as the time step. We denote by $t^n = n\Delta t$, and
 422 we assume that there exists $N \in \mathbb{N}^*$ such that $T = N\Delta t$. To obtain the fully discrete
 423 scheme associated with (29), we apply an explicit second-order time scheme centered
 424 at t^n , *a.k.a.* the leapfrog scheme [25, 28]. Namely, denoting by $\tilde{\mathbf{u}}_h^n$ the approximation
 425 of $\tilde{u}_h(t^n)$ solution of (29), we define the following fully discrete scheme:

426 For any $n \in \llbracket 1; N \rrbracket$, find $\tilde{\mathbf{u}}_h^{n+1} \in \mathcal{V}_h(\hat{\Omega})$ the solution, for any $\mathbf{w}_h \in \mathcal{V}_h(\hat{\Omega})$, of

$$\left(\frac{\tilde{\mathbf{u}}_h^{n+1} - 2\tilde{\mathbf{u}}_h^n + \tilde{\mathbf{u}}_h^{n-1}}{\Delta t^2}, \mathbf{w}_h \right)_h + a_h(\mathbf{u}_{0,H}; \tilde{\mathbf{u}}_h^n, \mathbf{w}_h) = \langle L(t^n, \mathbf{u}_{0,H}), \mathbf{w}_h \rangle. \quad (30)$$

427 with given initial steps $\tilde{\mathbf{u}}_h^0$ and $\tilde{\mathbf{u}}_h^1$.

428 From this expression of the fully discrete problem, it is possible to recover a discrete
 429 energy conservation property, which is the exact translation of (16). Denoting by $\|\cdot\|_h$
 430 the norm associated with the discrete inner-product $(\cdot, \cdot)_h$, we define the kinematic
 431 and potential discrete energies associated with the unique solution of (30) as

$$\mathcal{K}_h^{n+1/2} = \frac{1}{2} \left\| \frac{\tilde{\mathbf{u}}_h^{n+1} - \tilde{\mathbf{u}}_h^n}{\Delta t} \right\|_h^2, \quad \mathcal{P}_h^{n+1/2} = \frac{1}{2} a_h(\mathbf{u}_{0,H}; \tilde{\mathbf{u}}_h^{n+1}, \tilde{\mathbf{u}}_h^n),$$

432 and the total energy as $\mathcal{E}_h^{n+1/2} = \mathcal{K}_h^{n+1/2} + \mathcal{P}_h^{n+1/2}$. Assuming a null source term
 433 in (30), namely $L(t^n, \mathbf{u}_{0,H}) = 0$, then, taking $\mathbf{w}_h = \frac{1}{2\Delta t}(\tilde{\mathbf{u}}_h^{n+1} - \tilde{\mathbf{u}}_h^{n-1})$ as test function
 434 and using the symmetry of the stiffness bilinear form, we have the following energy
 435 conservation

$$\frac{\mathcal{E}_h^{n+1/2} - \mathcal{E}_h^{n-1/2}}{\Delta t} = 0. \quad (31)$$

436 This conservation property is critical since it leads to the conditions – on the stiffness
 437 bilinear form and on the time step – for the fully discrete (30) to be stable. Indeed,
 438 note that the total energy can be re-written in the following form

$$\begin{aligned} \mathcal{E}_h^{n+1/2} = \frac{1}{2} \left\| \frac{\tilde{\mathbf{u}}_h^{n+1} - \tilde{\mathbf{u}}_h^n}{\Delta t} \right\|_h^2 - \frac{\Delta t^2}{8} a_h \left(\mathbf{u}_{0,H}; \frac{\tilde{\mathbf{u}}_h^{n+1} - \tilde{\mathbf{u}}_h^n}{\Delta t}, \frac{\tilde{\mathbf{u}}_h^{n+1} - \tilde{\mathbf{u}}_h^n}{\Delta t} \right) \\ + a_h \left(\mathbf{u}_{0,H}; \frac{\tilde{\mathbf{u}}_h^{n+1} + \tilde{\mathbf{u}}_h^n}{2}, \frac{\tilde{\mathbf{u}}_h^{n+1} + \tilde{\mathbf{u}}_h^n}{2} \right). \end{aligned}$$

439 Hence, upon the assumption that

- 440 • The stiffness bilinear form in (30) satisfies the coercivity property (15),
- 441
- 442 • The time step satisfies

$$\Delta t \leq 2 \left(\sup_{\mathbf{w}_h \in \mathcal{V}_h(\hat{\Omega})} \frac{a_h(\mathbf{u}_{0,H}; \mathbf{w}_h, \mathbf{w}_h)}{\|\mathbf{w}_h\|_h^2} \right)^{-\frac{1}{2}}, \quad (32)$$

we can ensure that the total energy is a positive functional and that it actually
 corresponds to an equivalent norm on the discrete solution. Following arguments
 provided *e.g.* in [28, 43, 63] one can obtain the following stability result: there exists a
 constant $C > 0$ that depends on the initial steps $\tilde{\mathbf{u}}_h^0$ and $\tilde{\mathbf{u}}_h^1$, the source term $L(t^n; \mathbf{u}_{0,H})$
 and the maximal time T , such that the solution of (30) satisfies

$$\|\tilde{\mathbf{u}}_h^{n+1}\|_h \leq C(T, L, \tilde{\mathbf{u}}_h^0, \tilde{\mathbf{u}}_h^1), \quad \forall n \in \llbracket 1; N \rrbracket.$$

443 Remarks on practical considerations.

444 Prior to giving substantiated numerical results proving the soundness of the presented
445 modeling strategy, let us gather some important remarks.

446 First, in the definition of the discrete space $\mathcal{V}_h(\widehat{\Omega})$ we consider anisotropic orders
447 of approximation. This aspect appears to be particularly convenient for thin and
448 potentially stratified materials. Indeed, it allows for high-order elements in the mid-
449 surface of the domain – with low numerical dispersion or numerical anisotropy [25, 31,
450 34] – while maintaining a reasonable number of Degrees of Freedom (DoF) by using
451 lower order in the thickness of the material.

452 Second, the use of approximated numerical integration as denoted in (28) implies
453 in particular a lumped, *i.e.* diagonal, mass matrix. This is an important aspect in
454 the context of high-frequency wave propagation since the leapfrog scheme (30) thus
455 becomes fully explicit. At each time step, only a diagonal matrix needs to be inverted.
456 Additionally, using numerical integration with points matching the ones defining
457 the local Lagrange polynomial leads to very sparse discrete gradient computations.
458 Hence, the approximation $a(\mathbf{u}_0; \cdot, \cdot) \approx a_h(\mathbf{u}_0; \cdot, \cdot)$ also implies significant performance
459 gains – see *e.g.* [58] for more details on these practical aspects.

460 Moreover, in practice, the mesh step for the locking-free shell elements depends
461 mostly on the geometrical characteristics of the domain and the material variations –
462 typically one element per material layer in the case of stratified materials. For the wave
463 problem, to these constraints is added the constraint of having enough discretization
464 points per wavelength. With these qualitative meshing directives in mind, in the
465 context of high-frequency wave propagation, one can expect that the \mathcal{T}_h is significantly
466 more refined than \mathcal{T}_H , *i.e.* $h \ll H$.

In light of the previous remark, in order to solve the fully discrete one-way coupled
problems, one needs to resort to a mean of interpolating $\mathbf{u}_{0,H} \in \mathcal{V}_H(\widehat{\Omega})$ onto $\mathcal{V}_h(\widehat{\Omega})$.
Let us denote by $\mathcal{I}_h(\cdot)$ the nodal interpolator from $\mathcal{V}_H(\widehat{\Omega})$ to $\mathcal{V}_h(\widehat{\Omega})$, without entering
into the details of the exact expression of this interpolator, let us simply remark that
the stiffness bilinear form actually reads

$$a_h(\mathbf{u}_{0,H}; \cdot, \cdot) = a_h(\mathcal{I}_h(\mathbf{u}_{0,H}); \cdot, \cdot).$$

467 Finally, the condition (32) on the time step is often referred to as the Courant
468 – Friedrichs – Lewy (CFL) condition of the fully discrete scheme. Note that in the
469 particular cases of interest in our work, this CFL condition naturally depends on the
470 linearization point, *i.e.* the displacement field $\mathbf{u}_{0,H}$. Furthermore, in terms of stability,
471 in addition to this CFL condition, the coercivity assumption is capital – a condition
472 that is also required for the well-posedness of the continuous problem.

473 3 Numerical results and experimental validation

474 In this section, we apply our model and numerical methods for different configurations.
475 First, we assess the stability issue caused by a potential failure to ensure coercivity.
476 Second, we apply our methods to realistic cases in the SHM context. At last, we present
477 how we validate our model with experimental data available in the literature. The
478 computations were done using a laptop workstation equipped with a *Intel i9-9880H*
479 CPU and 32Gb of RAM. The wave propagation simulations are performed within the
480 SFEM kernel of the CIVA platform¹ developed at CEA.

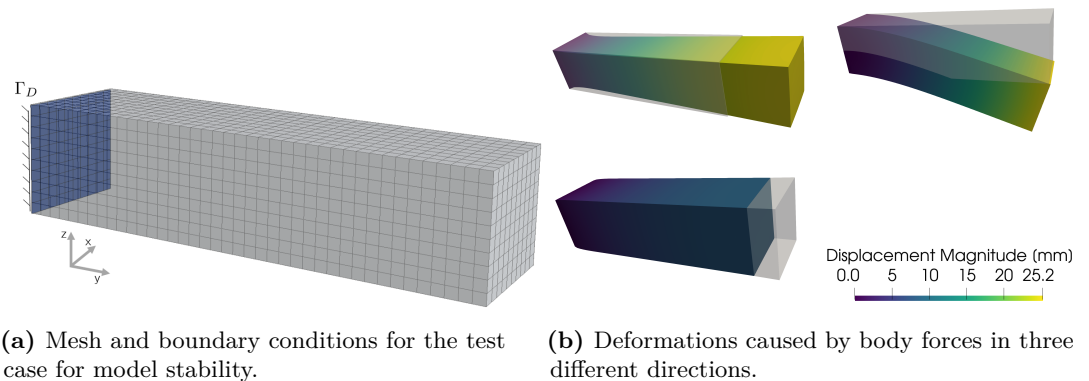
¹<https://www.extende.com>

Material	Density(g/cm ³)	λ (GPa)	μ (GPa)	l(GPa)	m(GPa)	n(GPa)
Aluminum	2.704	54.308	27.174	-181.0	-289.0	-336.0
Steel	7.850	111	82.1	-459.0	-461.0	-358.0

Table 1. Material parameters for the metals used in this work.

481 3.1 Illustration of potential stability issues of the fully discrete scheme

482 As discussed previously, the tangent stiffness operator in (14) must satisfy the coercivity
 483 assumption (15), otherwise the well-posedness of the problem cannot be guaranteed.
 484 This assumption cannot be satisfied for any deformation or constitutive law. In this
 485 section, we aim to illustrate this problem in a test model. At a discrete level, coercivity
 486 is equivalent to positiveness of the operator, which can be assessed by computing
 487 its eigenvalues. In the zero-deformation case, the presented model equates to linear
 488 elasticity (Hooke's law), and $a(0; \cdot, \cdot)$ is known to be coercive [17]. Here, we assess
 489 stability for deformed cases by analyzing the evolution of the system energy and the
 490 eigenvalues of the local stiffness matrices.



(a) Mesh and boundary conditions for the test

(b) Deformations caused by body forces in three different directions.

Figure 3. Configuration and deformations for illustrating the potential stability issues.

491 If instability would appear by the deformation, it could be interpreted as a shift
 492 of tangent stiffness operator eigenvalues towards negative values. In this case, the
 493 negative eigenvalues that are related to the instability of the scheme are expected
 494 to be among those of the smallest amplitude, *i.e.* low spatial frequency. Computing
 495 the lowest eigenvalues for the stiffness operator may have prohibitive computational
 496 cost due to the size of the problems at hand. To circumvent this computational
 497 difficulty, we analyze instead the eigenvalues of the local stiffness matrix related to
 498 each finite element. Although negative eigenvalues for the local stiffness matrices do
 499 not necessarily mean negative eigenvalues for the global matrix, it does give an insight
 500 into the problematic regions and related severity. In addition to the study of the lowest
 501 eigenvalue, we analyze the evolution in time of the total system's kinetic and potential
 502 energy for different initial conditions, showing the nature of such stability issues.

503 The model configuration used is an aluminum cuboid of 25mm \times 25mm \times 100mm
 504 depicted in *Figure 3a*. Dirichlet boundary condition is considered at $x=0$. The
 505 material parameters are present in *Table 1*. We analyze different constitutive laws
 506 and types of mechanical solicitations. The analysis is done using the aforementioned
 507 SVK, CNH, and Murnaghan's constitutive laws. For each case, three different types of
 508 mechanical solicitations as body forces will be considered:

- 509 • **Stretching:** 300N/mm³ (Y direction),
- 510 • **Bending:** 7.5N/mm³ (-Z direction),

511 • **Compressing:** $135\text{N}/\text{mm}^3$ (-Y direction).

512 The associated deformations are illustrated in *Figure 3b* for the CNH case. The
513 SVK law has slightly different deformation with the same nature and amplitudes.
514 Computing the quasi-static deformation using Murnaghan's constitutive law did not
515 converge due to its particularly ill-posed tangent operator.

516 After computing the deformations caused by these forces for each constitutive law,
517 the displacement field is given as input to our wave propagation solver. The same
518 constitutive law is used for the quasi-static problem and the wave propagation problem,
519 except in Murnaghan case. For the Murnaghan case, we compute the deformation
520 with CNH instead. The deformation field is interpolated from the quasi-static problem
521 mesh to the wave propagation problem mesh. After discretizing the functional space
522 using spectral finite elements, we assemble the local tangent stiffness matrices and
523 compute their eigenvalues. The minimum eigenvalues, for each element, are shown in
524 *Figure 4* for different types of deformations and constitutive laws.

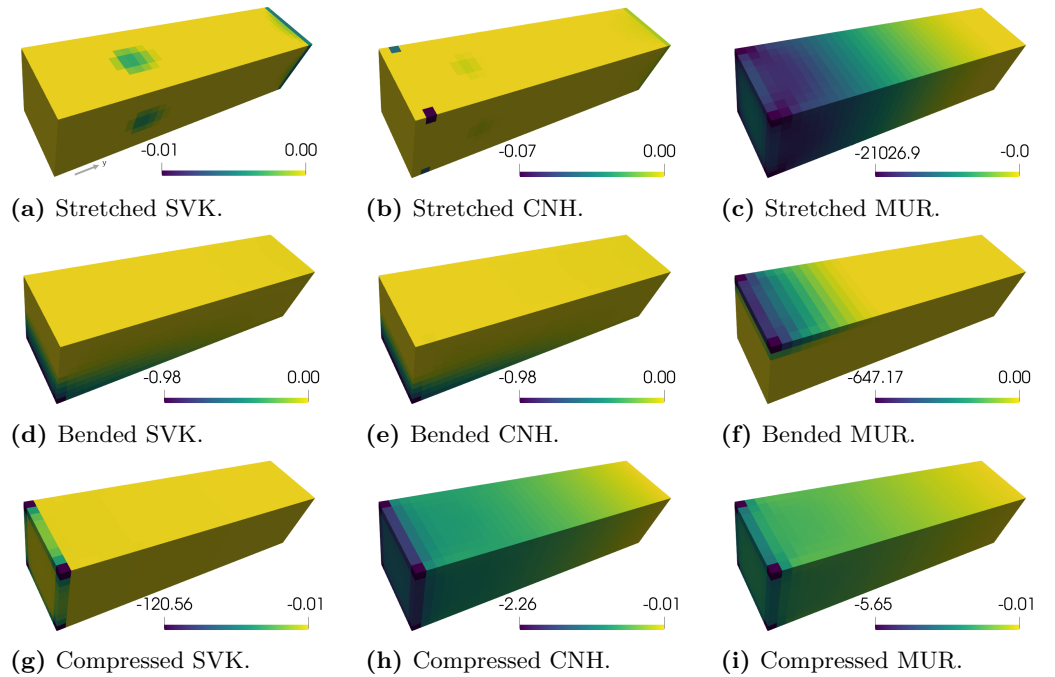
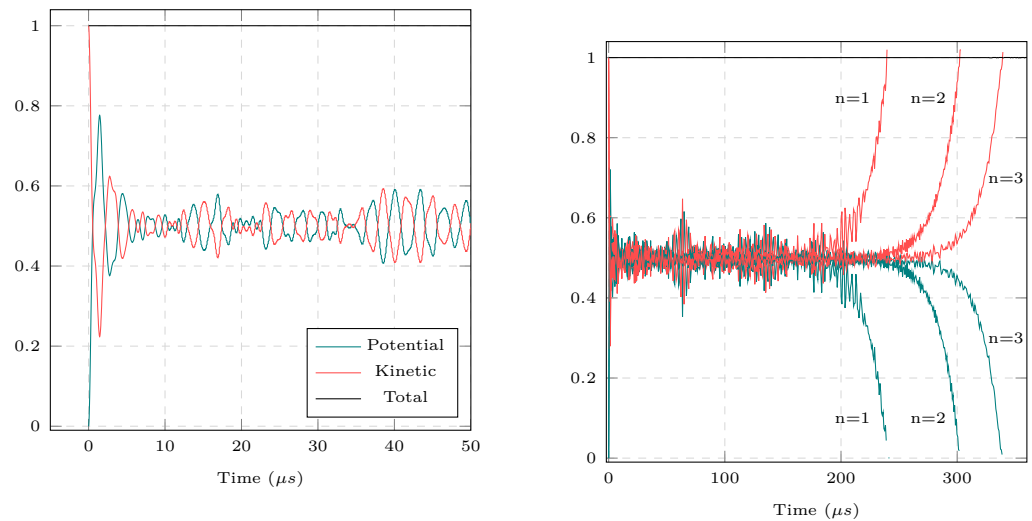


Figure 4. Minimum (local) tangent stiffness matrix eigenvalues for different types of deformation and constitutive laws.

525 The CNH constitutive law presents less severity, *i.e.* lowest amplitude negative
526 eigenvalues. The SVK and CNH cases are more unstable for compressed cases – cases
527 associated with buckling phenomena. The Murnaghan's law presents the highest
528 negative eigenvalues with critical behavior in the stretched case. The stability issue can
529 be also assessed by observing the evolution of the system's potential and kinetic energy.
530 We run the wave propagation problem, without source-term, from initial conditions
531 with specific spatial frequency signature, namely for $n = 1, 2, 3$

$$\begin{aligned}\tilde{\mathbf{u}}_h^0 &= - \left[\sin \left(\frac{25-x}{25} 2\pi n \right) \sin \left(\frac{25-y}{25} 2\pi n \right) \sin \left(\frac{25-z}{25} 2\pi n \right) \right] (\mathbf{e}_x + \mathbf{e}_y + \mathbf{e}_z), \\ \tilde{\mathbf{u}}_h^1 &= -\tilde{\mathbf{u}}_h^0.\end{aligned}$$

532 It allows us to solicit the low-frequency modes and assess the instability, relating the
533 frequency content to the issue. Running simulations with these initial conditions, we
534 monitor the total, kinetic and potential energies as shown in *Figure 5*.



(a) Stable behavior for a stable case (Stretched; CNH). (b) Unstable case (Bended; CNH) for different initial conditions.

Figure 5. Evolution of total energies (*a.u.*) for the wave propagation problem.

535 As the total energy is conserved, if both kinetic and potential energies are positive,
 536 the system's energy is bounded. However, if any of the energies become negative, the
 537 solution may diverge. For the stable cases, the energy transfer between kinetic and
 538 potential is kept throughout the simulation. For the unstable cases, as soon as the
 539 potential energy presents negative values, the system starts to diverge. The excitation
 540 of lower frequencies presented instability in fewer time steps as they excite lower
 541 frequency modes, *i.e.* modes related to lower eigenvalues. For the presented forces,
 542 the CNH case is unstable only under compression, SVK is unstable for bending and
 543 compression and Murnaghan's is unstable for all presented deformations. In general,
 544 the stability issues are not only dependent of buckling phenomena but also on the
 545 choice of constitutive law.

546 3.2 Illustrations related to *Structural Health Monitoring*

547 We now illustrate the generic aspect of our work by applying it to different cases.
 548 The first case is a steel pipe subjected to the 4-point bending test, inspired by an
 549 experiment done in [54]. We simulate ultrasonic propagation with different deformation
 550 intensities and compare the acoustoelastic effects in the signal with ones caused by
 551 a crack through the thickness of the pipe. In the second case, we model ultrasonic
 552 wave propagation in a 16-layer stratified composite plate under torsion forces. For
 553 such modeling, we use the transversely isotropic law (22) for each of the 16 layers
 554 associated with a principal fiber direction. These examples illustrate the application of
 555 our method to realistic and complex loading configurations, and the simulation results
 556 could be used to assess the efficiency of a SHM system.

557 Illustration on a pipe subject to a 4-point bending test.

558 We present here a case of wave propagation in a steel pipe subjected to a 4-point
 559 bending test. This configuration is usually used to perform fatigue tests in welded
 560 pipes. The experimental setup is inspired by the experiment done in [54]. We model
 561 and simulate the wave propagation in the pipe subjected to such bending forces
 562 with different intensity levels. We use Murnaghan's constitutive law with material
 563 parameters for steel (*Table 1*). For these levels of deformation and material parameters
 564 the wave propagation problem does not present stability issues. The surface force is

565 applied at two regions close to the pipe extremities. The configuration is depicted
 566 in *Figure 6*, with a 2.94m long, 8mm thick pipe and an inner radius of 102.65mm.
 567 The resulting deformation for a total force of 220kN is shown in *Figure 7*.

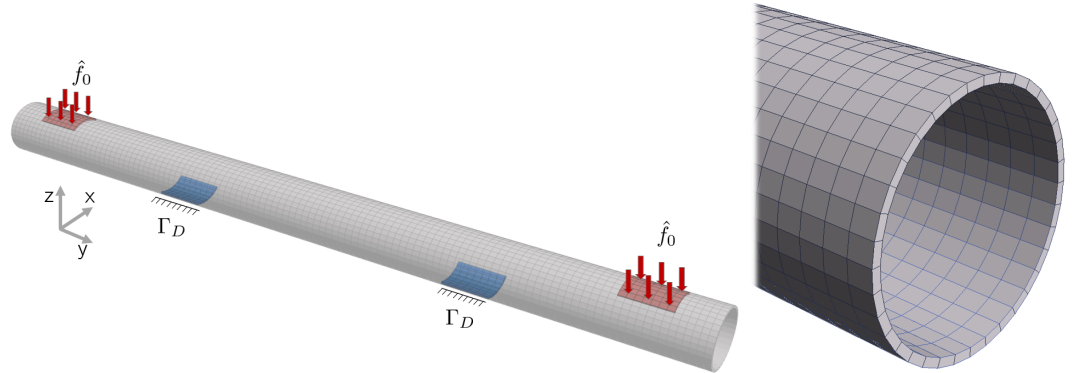


Figure 6. Numerical setup and mesh for the quasi-static problem.

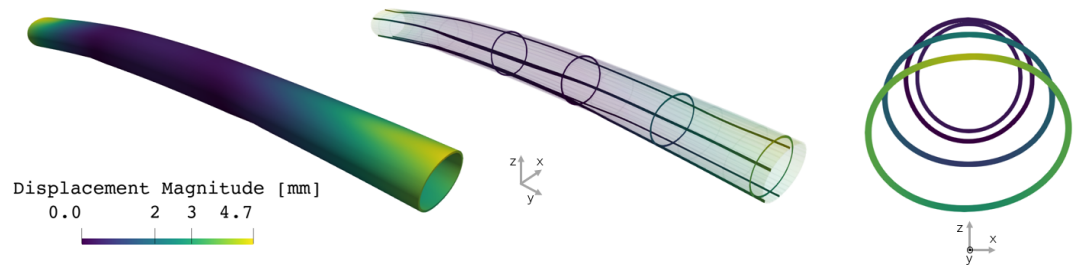


Figure 7. Deformation of the pipe for a 220kN loading. (Deformation scaled 20x for visualization).

568 Given the different deformations as input, we interpolate them in the wave propaga-
 569 tion problem mesh and compute the wavefield generated by a unique source term. The
 570 wave problem mesh, source location and measurement point are depicted in *Figure 8*,
 571 as well as the position of a potential crack. The ultrasonic waves are excited close
 572 to one edge of the pipe and measured at the other edge. For a case with no cracks,
 573 we simulate the wavefield and extract the signal from the modeled transducer at the
 574 measurement point and plot in *Figure 9* for different levels of mechanical loadings.
 575 Snapshots of the wavefield are shown in *Figure 10*.

576 In *Figure 9*, we can see that the effect of bending forces results in a change in
 577 time of flight, decreasing the wave velocity, and also a change in amplitude of the
 578 signals with higher forces.

579 To evaluate how the effect of mechanical loading on the signal compares with the
 580 one caused by a crack, we set up and simulate three numerical configurations:

- 581 • S : non deformed pipe (considered as baseline; no defects)
 582 *Simulation is done considering no deformation with a pristine pipe,*
- 583 • S_p : deformed pipe (without defect)
 584 *Simulation is done considering the pipe under bending forces of 220kN,*
- 585 • S_d : pipe with defect (without deformation)
 586 *Simulation is done considering the undeformed pipe with a crack of 19mm wide*
 587 *through half the thickness.*

588 From each case, we extract the signal of a modeled transducer at the measurement point.
 589 The defect through half the thickness was modeled by disconnecting elements. By

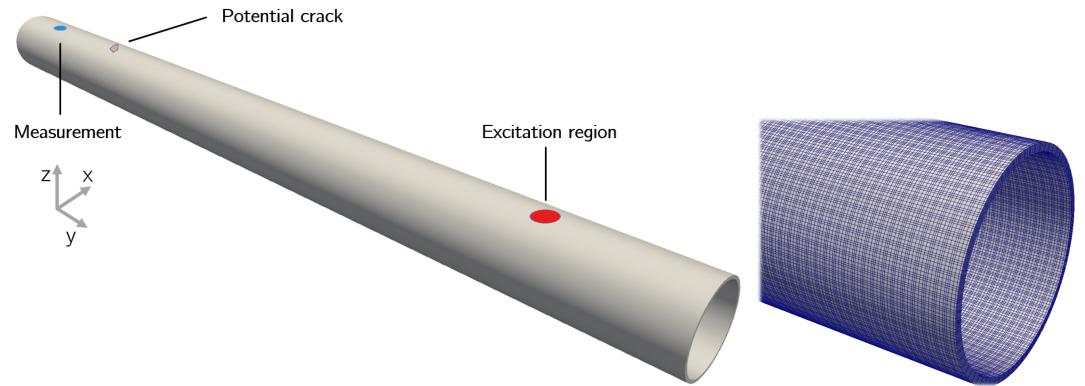


Figure 8. Transducer, potential crack and measurement point positions for the wave propagation problem. On the right-hand side, the mesh for the wave propagation problem characterized by the Gauss-Lobatto points.

Material	Density(g/cm ³)	C_{11}	C_{22}	C_{12}	C_{23}	C_{66}
CFRP	1.6	155.17	13.22	5.34	6.98	5.77

Table 2. Material parameters for CFRP. [29]

590 subtracting the signals from the baseline, we can isolate the individual contribution to
 591 the signal of the mechanical loading and the crack. The baseline signal S is subtracted
 592 from S_p and S_d and the results are plotted in *Figure 11*.

593 Regarding performances, the quasi-static problem is 148k DoFs and takes 3 minutes
 594 to run 5 Newton-Raphson iterations while using 2.2Gb of RAM. The wave propagation
 595 problem requires around 3.86M mesh points to ensure at least two elements of fourth-
 596 order per wavelength, resulting in 11.58M DoFs. The timestep is $\Delta t = 0.11\mu\text{s}$ to
 597 comply with the CFL condition (32). Using an efficient implementation of the spectral
 598 finite elements, the numerical solver uses 2.3Gb of RAM, 73 seconds for initialization
 599 (starting the model, initial outputs and interpolating the structure displacement field)
 600 and 40 minutes to compute approximately 4336 timesteps (including wavefield outputs).

601 The effect on the signal caused by the deformation alone is comparable to the
 602 effect caused by a crack for such configuration and type of measurement. Thus,
 603 not considering these effects when analysing wave propagation data for this level of
 604 mechanical loading results in less reliable detection and evaluation of such type of
 605 damages.

606 Illustration on a stratified anisotropic plate.

607 Carbon fiber-reinforced polymers are largely employed in aerospace industries, where
 608 ensuring safety is essential. They are made of a polymer matrix reinforced with
 609 stacked layers of fibers. Usually, each fiber layer has a principal direction, presenting
 610 a transversely isotropic response to mechanical solicitations. Here, we model the
 611 wave propagation in a 16-layer Carbon Fiber Reinforced Polymer (CFRP) plate with
 612 dimensions of $900 \times 300 \times 2$ (16×0.125) mm³ under torsion as illustrated in *Figure 6*.

Each layer is modeled using the hyperelastic potential (22) where the direction related to the invariants I_4 and I_5 were set accordingly to the fiber directions in *Figure 12*. The material parameters are shown in *Table 2*. The deformation was caused by setting the Dirichlet boundary condition at $x = 0$ and applying a body force with the following form

$$f = \frac{y - 150}{300} 10^{-3} \text{N/mm}^3, \quad \forall y \in [0, 300].$$

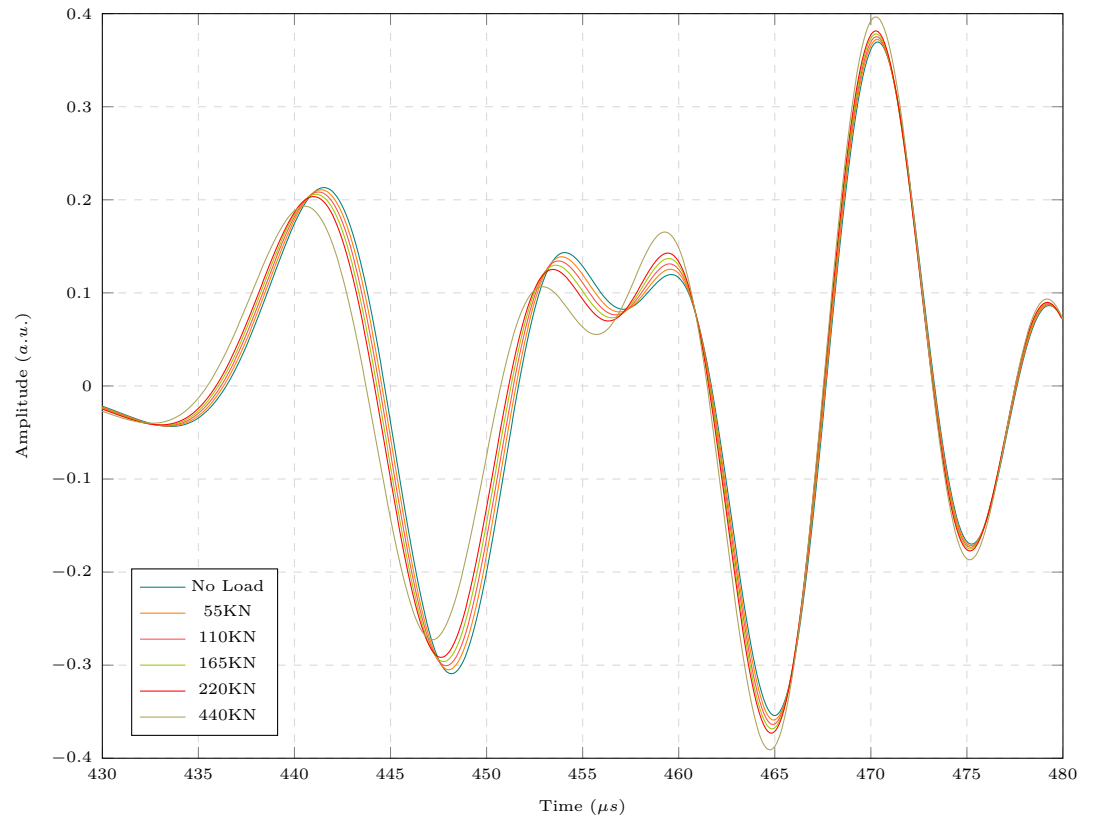


Figure 9. Transducer signal at the measurement point for different bending forces.

613 The resulting deformation by solving the quasi-static problem is shown in *Figure 13*.
 614 Snapshots of the wavefield computed solving the wave propagation problem can be
 615 seen in *Figure 14*.

616 The quasi-static problem has 270k DoFs and takes 50 minutes to run 10 Newton-
 617 Raphson iterations while using 12Gb of RAM. The wave propagation problem requires
 618 around 8.3M mesh points, having at least two elements of fourth-order per wave-
 619 length, resulting in 25M DoFs. Satisfying the CFL condition (32), the timestep is
 620 $\Delta t = 0.0159\mu s$. The wave propagation solver requires 4Gb of RAM, 9 minutes for
 621 initialization (starting the model, initial outputs and interpolating the structure dis-
 622 placement field) and 2 hours and 54 minutes to compute approximately 18867 timesteps
 623 (including wavefield outputs). This computation is particularly expensive as we need
 624 to satisfy the CFL condition for the mesh step related to the thickness discretization
 625 (16 elements for 2mm). This problem could be mitigated by using implicit-explicit
 626 schemes as in [64].

627 This illustration shows the capability of our simulation tools to model thin structures
 628 with the possibility of choosing a hyperelastic law that best suits the modeling objectives.
 629 As discussed in the paragraph about acoustoelasticity, the acoustoelastic effect is
 630 controlled by the TOECs. Here, we choose to use a hyperelastic law that matches only
 631 second-order elastic constants although the implementation of laws that include the
 632 TOECs can also be considered.

633 3.3 Experimental validation on an isotropic plate

634 Finally, in this section, we model and simulate an experimental setup to validate our
 635 model in the isotropic case. The experiment is presented in [41], where the authors

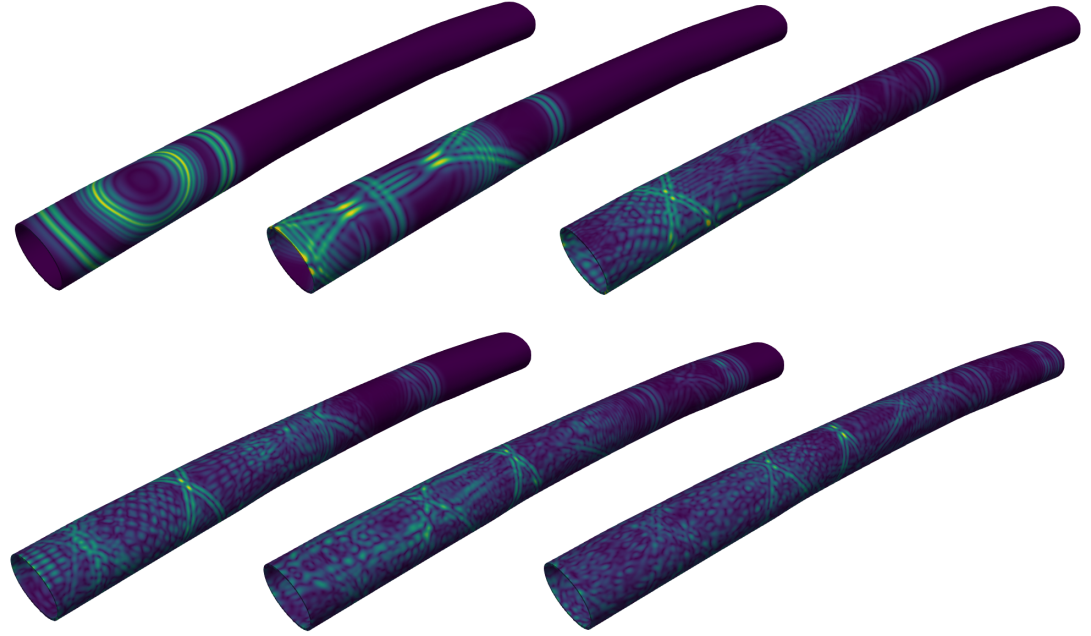


Figure 10. Snapshots of the wavefield computed in the steel pipe subjected to the 4-point bending test at $t = 70, 140, 230, 300, 390$ and $480\mu\text{s}$ (left to right, top to bottom).

Material	Density (g/cm^3)	λ (GPa)	μ	l	m	n
Aluminum (fitted)	2.704	54.308	27.174	-181.0	-289.0	-336.0

Table 3. Material parameters for aluminum, fitted to experimental data.

636 apply axial forces to an aluminum plate and measure the change in wave speed caused
 637 by the load-induced deformation, for different directions and guided wave modes.
 638 Using Murnaghan’s constitutive law, the authors propose an eigenproblem to obtain
 639 guided wave modes and their wave speeds in biaxially stressed isotropic materials.
 640 By applying canonical deformations and measuring the change in wave speed for
 641 different angles, the authors calibrated the TOECs of Murnaghan’s constitutive law
 642 using experimental data. The obtained parameters are presented in *Table 3*. Here, we
 643 aim at modeling and simulating this experiment using the hyperelastic potential (21)
 644 with these parameters. We compare the simulated results with their experimental data
 645 to validate our model and numerical tools. The experiment is depicted in *Figure 15*.

646 In the experimental setup, eleven levels of axial loading are applied from 0MPa
 647 to 57.5MPa and the ultrasonic wave speeds are measured for 9 different angles of
 648 propagation with transducers positioned at 109mm from the center in the undeformed
 649 configuration [37]. For the validation procedures, we use their measured changes in
 650 wave speed for the S0 mode excited with a 250kHz 5-cycle Hanning windowed cosine.
 651 The wave speed is obtained by retrieving the ToF between transducer pairs for the
 652 S0 mode at a given load. To avoid effects of dispersion on the ToF measurement, we
 653 identify the 4th zero-crossing of the signal as the least affected by dispersion. As our
 654 solution is computed in the reference configuration, there is no need for correcting the
 655 distance change between transducers caused by the deformation. The authors in [41]
 656 fitted their parameters while neglecting this distance difference when computing the
 657 wave speed using the ToF. The computed displacement between the axial pair (#1,#6)

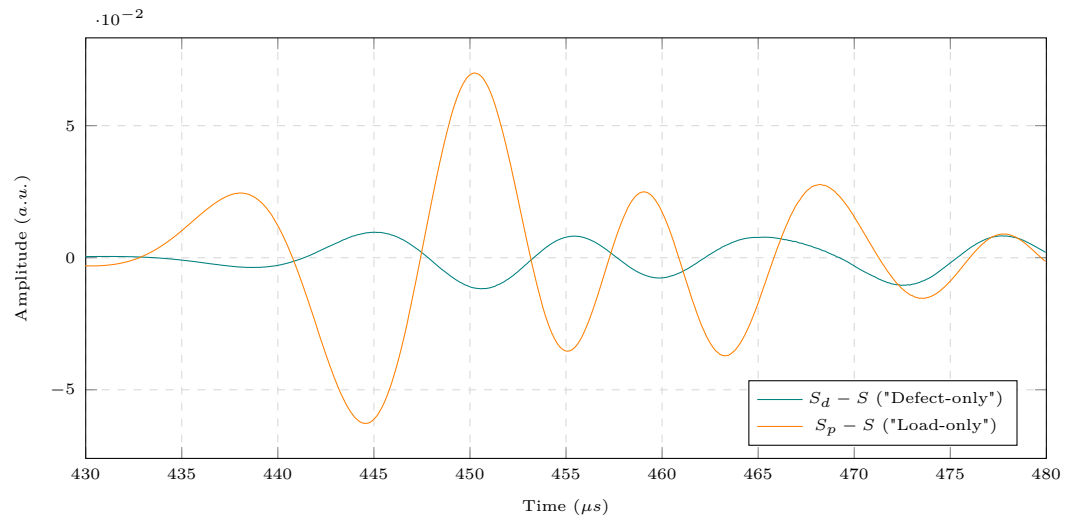


Figure 11. Transducer signal at measure point subtracted from the baseline for the case with defect-only and load-only.

658 has a maximum of 0.173mm at 57.5MPa. It represents a change in the total pair
 659 distance of around 0.08%, or, in terms of wave speed, around 4.2m/s when considering
 660 the mean wave speed of 5300m/s.

661 We simulate and obtain, for the different load levels, the changes in time of flight
 662 (ToF) for the angles 90° , 45° , 0° , -22.5° , -67.5° , corresponding to the transducer pairs
 663 ($\#1, \#6$), ($\#5, \#10$), ($\#4, \#9$), ($\#2, \#7$) and ($\#3, \#8$). The Y axis is considered to be
 664 at an angle of 90° . The changes in wave speed in the loading axis with respect to
 665 the load intensity are plotted in *Figure 16a*. Here, simulated results with different
 666 constitutive laws (SVK, CNH and Murnaghan's) are compared to experimental data
 667 and the analytical reference solution proposed in [41]. The behavior of SVK and CNH
 668 constitutive laws reflect the analysis done in the paragraph about acoustoelastic effects.
 669 With the SVK law, increasing axial load leads to increasing incremental wave speeds,
 670 what is opposite to the observed effect. Although the CNH models the right tendency of
 671 decreasing the wave speed, it underestimate its effects. Making use of the TOECs, the
 672 Murnaghan's law is able to properly model the experimental data. For the maximum
 673 load of 57.5MPa, the changes in wave speed with respect to the angle of propagation
 674 are plotted in *Figure 16b*. Here, our method implemented with the Murnaghan's
 675 law is also compared with the experimental data and the analytical proposition when
 676 obtaining the angular-dependency of the changes in wave speed. With our simulation,
 677 we retrieve the expected bell-shaped curve, effect of the load-induced anisotropy also
 678 obtained from the ToF measurements from experimental and simulated data.

679 Summarizing our validation achievement, the experiment is modeled and simulated
 680 starting from the mechanical loading that results in the deformation of the aluminum
 681 plate. This deformation is used as input to our linearized wave propagation model,
 682 and, using the proposed constitutive law and material parameters we are able to
 683 extract the ultrasonic signals of the same nature as in the experiment. Our simulation
 684 reliably matches the experimental results while computing the whole wave field. We
 685 also implement different constitutive laws, obtaining their expected behavior. This
 686 further validates each modeling component and solvers used.

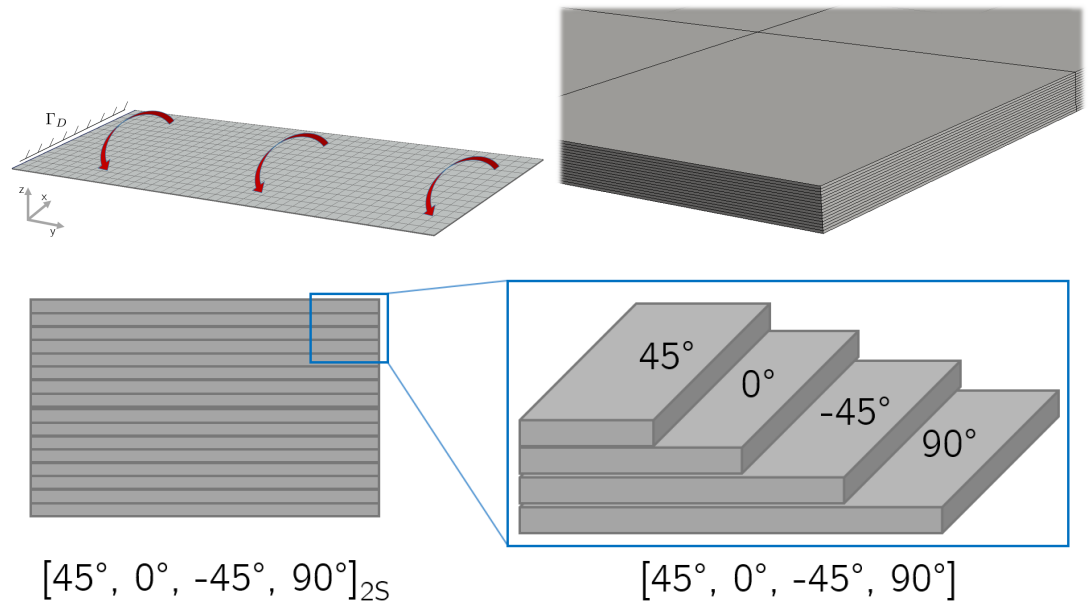


Figure 12. Configuration for the quasi-static problem (upper left). Zoom showing the stratified aspect (upper right). Ply configuration (bottom).

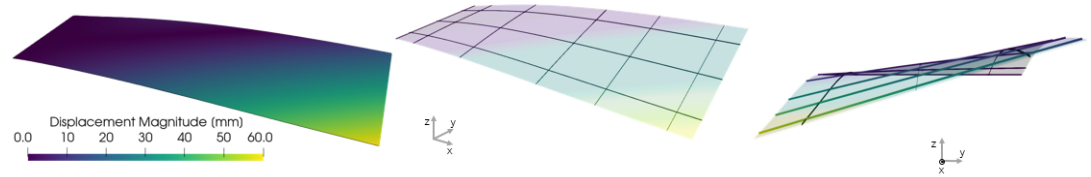


Figure 13. Deformation of the CFRP plate under torsion. (visualization scaled 2x).

687 4 Conclusion & perspectives

688 In our work, we have proposed a set of numerical tools addressing the problem of
 689 ultrasonic guided-wave propagation in materials subject to pre-deformation. Modeling
 690 of such configurations includes two different problems: a quasi-static (time-independent)
 691 nonlinear problem satisfied by the large displacement field, and a linearized time-domain
 692 wave propagation problem. Our approach is generic in the sense that it can represent
 693 any type of constitutive law and geometries. For each problem, we have carefully
 694 chosen a dedicated numerical method – namely the 3D shell elements for the first
 695 problem and the time-domain SFEM for the second one – leading to computationally
 696 efficient algorithms. We have presented numerical results on 3D cases related to SHM
 697 configurations, thus illustrating the performances of our method for both isotropic
 698 and anisotropic materials. We also have compared our results against experimental
 699 data extracted from the literature, thus validating the overall modeling and numerical
 700 toolchain.

701 Thanks to the presented approach, we were able to illustrate how the effects of
 702 a pre-deformation can alter the ultrasonic signals used to monitor a specimen under
 703 inspection. It makes this work in a good position to get closer to real applications in
 704 industrial contexts, and to provide meaningful modeling tools for potential industrial
 705 partners aiming at designing SHM systems. In order to enhance the attractiveness
 706 of these modeling tools, it would be interesting to include a mean of automatically
 707 removing the bias induced by the pre-deformation or the EOCs in general. This

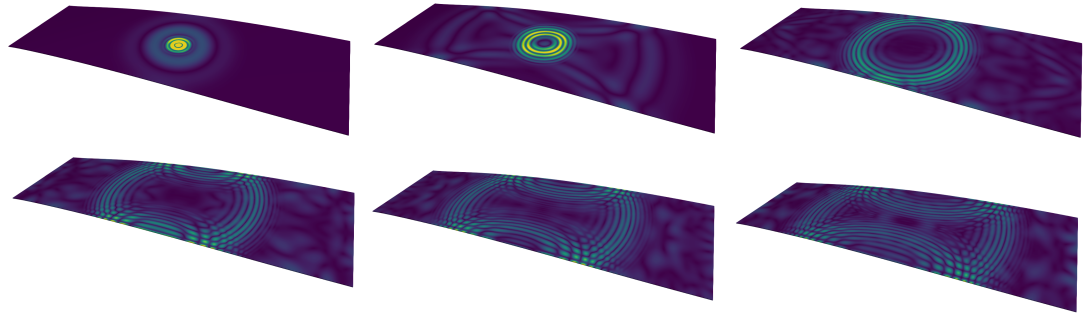


Figure 14. Wavefield computed in a 16 layers CFRP plate under torsion at $t = 45, 90, 180, 240, 270$ and $300\mu s$.

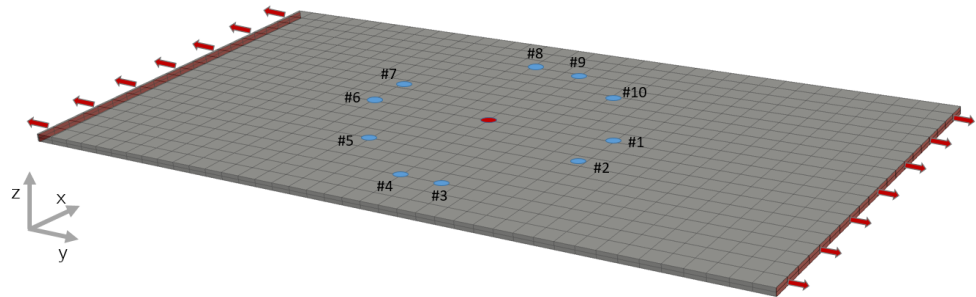
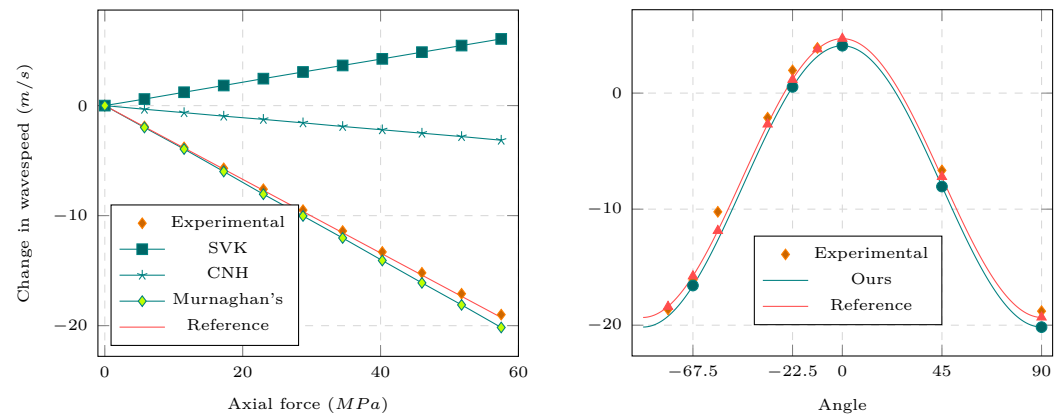


Figure 15. Experimental setup for calibrating the TOECs using the acoustoelastic effect.

708 could typically be done by considering an inverse problem where we would aim at
 709 reconstructing the pre-deformation sustained by the structure from the guided-wave
 710 measurements.

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(a) Load-induced wavespeed changes for the S0 guided wave mode at 250kHz. (b) Load-induced wavespeed changes for the S0 mode at 250kHz. (Murnaghan's law)

Figure 16. Comparison between simulated data (for different constitutive laws) and experimental data, validating our proposed model.

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