Theory of pressure acoustics with viscous boundary layers and streaming in curved elastic cavities

Bach, Jacob Søberg; Bruus, Henrik

Published in:
Journal of the Acoustical Society of America

Link to article, DOI:
10.1121/1.5049579

Publication date:
2018

Document Version
Publisher's PDF, also known as Version of record

Link back to DTU Orbit

Citation (APA):

General rights
Copyright and moral rights for the publications made accessible in the public portal are retained by the authors and/or other copyright owners and it is a condition of accessing publications that users recognise and abide by the legal requirements associated with these rights.

- Users may download and print one copy of any publication from the public portal for the purpose of private study or research.
- You may not further distribute the material or use it for any profit-making activity or commercial gain
- You may freely distribute the URL identifying the publication in the public portal

If you believe that this document breaches copyright please contact us providing details, and we will remove access to the work immediately and investigate your claim.
Theory of pressure acoustics with viscous boundary layers and streaming in curved elastic cavities

Jacob S. Bach and Henrik Bruus

Department of Physics, Technical University of Denmark, DTU Physics Building 309, DK-2800 Kongens Lyngby, Denmark

(Received 22 April 2018; revised 6 July 2018; accepted 20 July 2018; published online 17 August 2018)

The acoustic fields and streaming in a confined fluid depend strongly on the viscous boundary layer forming near the wall. The width of this layer is typically much smaller than the bulk length scale set by the geometry or the acoustic wavelength, which makes direct numerical simulations challenging. Based on this separation in length scales, the classical theory of pressure acoustics is extended by deriving a boundary condition for the acoustic pressure that takes viscous boundary-layer effects fully into account. Using the same length-scale separation for the steady second-order streaming, and combining it with time-averaged short-range products of first-order fields, the usual limiting-velocity theory is replaced with an analytical slip-velocity condition on the long-range streaming field at the wall. The derived boundary conditions are valid for oscillating cavities of arbitrary shape and wall motion, as long as both the wall curvature and displacement amplitude are sufficiently small. Finally, the theory is validated by comparison with direct numerical simulation in two examples of two-dimensional water-filled cavities: The well-studied rectangular cavity with prescribed wall actuation, and a more generic elliptical cavity embedded in an externally actuated rectangular elastic glass block. © 2018 Acoustical Society of America.

https://doi.org/10.1121/1.5049579

[PLM] Pages: 766–784

I. INTRODUCTION

The study of ultrasound effects in fluids in sub-millimeter cavities and channels has intensified the past decade, as microscale acoustofluidic devices are used increasingly in biology, environmental and forensic sciences, and clinical diagnostics. Examples include cell synchronization, enrichment of prostate cancer cells in blood, size-independent sorting of cells, manipulation of C. elegans, and single-cell patterning. Acoustics can also be used for non-contact microfluidic trapping and particle enrichment as well as acoustic tweezing.

The two fundamental physical phenomena that enable these microscale acoustofluidic applications are rooted in nonlinear acoustics. One fundamental phenomenon is the acoustic radiation force, which tends to focus suspended particles in the pressure nodes based on their acoustic contrast to the surrounding fluid. The second fundamental phenomenon is the acoustic streaming appearing as steady flow rolls which tend to defocus suspended particles due to the Stokes drag. Because the acoustic radiation force scales with the volume of the suspended particle, and the Stokes drag with its radius, the former dominates for large particles and the latter for small. For water at room temperature and 1 MHz ultrasound, the critical particle radius for the crossover between these two regimes has been determined to be around $2\,\mu m$.

So far, the vast majority of successful microscale acoustofluidics applications has been for large (above 2 $\mu m$) particles, such as cells, whose dynamics is dominated by the well-characterized, robust acoustic radiation force, which depends on the bulk properties of the acoustic field and material parameters of the particles and the surrounding fluid. However, there is a strong motivation to handle also sub-micrometer particles such as bacteria, exosomes, and viruses, for use in contemporary lab-on-a-chip-based diagnostics and biomedical research. In contrast to large particles, the dynamics of small (sub-micrometer) particles is dominated by the drag force from the ill-characterized acoustic streaming. To control the handling of such nanoparticle suspensions, a deeper understanding of the often complicated acoustic streaming is called for.

One important aspect of ultrasound acoustics is the large velocity gradients in the sub-micrometer-thin viscous boundary layer near rigid boundaries. The shear stress and the Reynolds stress that build up in this region are responsible for the viscous damping of the acoustic fields and for the acoustic streaming, respectively. In water with kinematic viscosity $\nu_0 \approx 10^{-6} \, m^2/s$ at the frequency $f = (1/2\pi)\omega \approx 1 \, MHz$, the thickness $\delta = \sqrt{2\nu_0/\omega}$ of this boundary layer is on the order of 500 nm, while the acoustic wavelength is around 1.5 mm. This three-orders-of-magnitude separation of physically relevant length scales poses a severe challenge for numerical simulations. To circumvent the problem of resolving the thin boundary layer, we develop a theory were analytical solutions for the boundary layers are used to formulate boundary conditions for the pressure field and bulk streaming field, which both varies on the much longer length scale $d \gg \delta$.

First, we extend the classical pressure acoustics theory by formulating a boundary condition for the acoustic pressure that includes the presence of the boundary layer, which

$^a$Electronic mail: bruus@fysik.dtu.dk
is otherwise neglected. Thus, our extended boundary condition takes into account important effect of the boundary layer, such as increased viscous damping, shifts in resonance frequencies, and shear stresses on the surrounding walls.

Second, we formulate a generalized slip-velocity boundary condition for bulk acoustic streaming over curved oscillating surfaces. An important step in this direction was the development of the limiting-velocity theory by Nyborg in 1958 for perpendicularly oscillating curved walls. Later modifications of this theory comprise modifications to the analysis in curvilinear coordinates by Lee and Wang in 1989, and the treatment of oscillations in any direction for flat walls by Vanneste and Bühler in 2011. Here, we extend these theories to harmonic oscillations in any direction of an arbitrarily shaped elastic wall, provided that both the radius of curvature and the acoustic wavelength are much larger than the boundary layer length-scale \( \delta \), and that also the amplitude of the perpendicular surface vibration is much smaller than \( \delta \).

Notably, the theoretical description developed here allows us to perform numerical simulations of the linear and nonlinear acoustics in arbitrarily shaped liquid-filled cavities, including the development of the limiting-velocity theory by Nyborg in 1958. Here, we extend this theory to account for important effects of the boundary layer and the acoustic wavelength, and shear stresses on the surrounding walls.

The instantaneous position \( s(s_0, t) \) at time \( t \) of this interface (the wall), is described by the small complex displacement \( s_1(s_0)e^{-i\omega t} \).

\[
s(s_0, t) = s_0 + s_1(s_0) e^{-i\omega t}. \tag{1}
\]

In contrast to Muller and Bruus, we do not study the transient phase leading to this steady oscillatory motion.

**A. Fundamental conservation laws in acoustofluidics**

The theory of acoustofluidics in \( \Omega \) is derived from the conservation of the fluid mass and momentum density,

\[
\begin{align*}
\partial_t \rho &= -\nabla \cdot (\rho \mathbf{v}), \\
\partial_t (\rho \mathbf{v}) &= -\nabla \cdot [(\rho \mathbf{v}) \mathbf{v}] + \nabla \cdot \mathbf{\sigma},
\end{align*}
\]

where \( \rho \) is the mass density, \( \mathbf{v} \) is the Eulerian fluid velocity, and \( \mathbf{\sigma} \) is the viscous stress tensor, given by

\[
\mathbf{\sigma} = -p \mathbf{I} + \mathbf{\tau},
\]

\[
\mathbf{\tau} = \eta_0 [\mathbf{v} \cdot \mathbf{v}] + \eta_1 \left[ \mathbf{v} + (\mathbf{v})^T - \frac{2}{3} (\mathbf{v} \cdot \mathbf{v}) \mathbf{I} \right].
\]

Here, \( p \) is the pressure and \( \mathbf{\tau} \) is the viscous part of the stress tensor given in terms of the bulk viscosity \( \eta_0 \), the dynamic viscosity \( \eta_1 \), the identity matrix \( \mathbf{I} \), and the superscript “\( T \)” denoting matrix transpose. Thermal dissipation is neglected throughout this work. We introduce the isentropic compressibility \( k_0 \) and speed of sound \( c_0 \),

\[
k_0 = \frac{1}{\rho_0} \frac{\partial \rho}{\partial p} = \frac{1}{\rho_0 c_0^2},
\]

as well as the small dimensionless damping coefficient \( \Gamma \) in terms of the viscosity ratio \( \beta \),

\[
\Gamma = (\beta + 1) \eta_0 k_0, \quad \beta = \frac{\eta_0}{\eta_1} + \frac{1}{3}.
\]

**B. Perturbation expansion**

The linear acoustic response of the system is proportional to the displacement stimulus \( s_1(s_0)e^{-i\omega t} \), and the resulting complex-valued quantities \( Q_1(r)e^{-i\omega t} \) are called first-order fields with subscript “1”. The physical time-dependent quantity \( Q_1^{phys}(r, t) \) corresponding to \( Q_1 \) is given by the real part \( Q_1^{phys}(r, t) = \text{Re}[Q_1(r)e^{-i\omega t}] \).

As the governing equations are nonlinear, we also encounter higher-order terms, and in the present work, we include terms to second order in the stimulus. Moreover, since we are only interested in the steady part of these second-order fields, we let in the following the subscript “2” denote a time-averaged quantity, written as \( \langle Q_2(r) \rangle = \langle Q_2(r, t) \rangle = \langle (\omega/2\pi)^{-2} Q_2(r, t) \rangle \rangle \). Time-averages of products of time-harmonic complex-valued first-order fields \( A_1 \) and \( B_1 \) are also of second order, and for those we have \( \langle A_1 B_1 \rangle = \frac{1}{2} \text{Re}[A_1(r)B_1^*(r)] \rangle \), where the asterisk denote complex conjugation.

Using this notation for the fluid, we expand the mass density \( \rho \), the pressure \( p \), and the velocity \( \mathbf{v} \) in perturbation series of the form

\[
\begin{align*}
\rho &= \rho_0 + \rho_1(r)e^{-i\omega t} + \rho_2(r), \\
p &= p_0 + p_1(r)e^{-i\omega t} + p_2(r), \\
\mathbf{v} &= \mathbf{v}_0 + \mathbf{v}_1(r)e^{-i\omega t} + \mathbf{v}_2(r),
\end{align*}
\]

**II. WALL MOTION AND PERTURBATION THEORY**

We consider a fluid domain \( \Omega \) bounded by an elastic, oscillating solid, see Fig. 1. All acoustic effects in the fluid are generated by the fluid-solid interface that oscillates harmonically around its equilibrium position, denoted \( s_0 \) or \( \partial \Omega \), with an angular frequency \( \omega \). The instantaneous position \( s(s_0, t) \) at time \( t \) of this interface (the wall), is described by the small complex displacement \( s_1(s_0)e^{-i\omega t} \).

\[
s(s_0, t) = s_0 + s_1(s_0) e^{-i\omega t}. \tag{1}
\]

**FIG. 1.** (Color online) Sketch of the interface between a fluid (light blue, \( \Omega \)) and a curved, oscillating solid (dark gray) with instantaneous position \( s_1(s_0) \) and equilibrium position \( s_0 \) (black line, \( \partial \Omega \)). The local curvilinear coordinate system on the interface is given by the tangent vectors \( \mathbf{e}_t \) and \( \mathbf{e}_n \) and the normal vector \( \mathbf{e}_\zeta \). By a Helmholtz decomposition, the first-order acoustic fluid velocity \( \mathbf{v}_1 = \mathbf{v}_1^d + \mathbf{v}_1^s \) is written as the sum of a long-range compressible part \( \mathbf{v}_1^d \) (dark blue) extending into the bulk and a short-range incompressible part \( \mathbf{v}_1^s \) (light red) with a decay length equal to the boundary-layer width \( \delta \). \( \mathbf{v}_1^0 = \mathbf{v}_1^d + \mathbf{v}_1^s \) is the Lagrangian velocity of the interface (the wall).
where $\rho_1 \ll \rho_0$, $\rho_1 = \frac{c_0^2}{\rho_0}$, and $|v_1| \ll c_0$. The subscripts 1 and 2 denote the order in the small acoustic Mach number $M_a = (1/c_0)|v_1|$, which itself is proportional to $s_1$.

### C. No-slip boundary condition at the wall

To characterize the wall motion, we compute the time derivative of $\mathbf{s}(s_0, t)$ in Eq. (1),

$$\partial_s s_0(t) = -i\omega \mathbf{s}_1(s_0) e^{-i\omega t} = V_0(s_0) e^{-i\omega t},$$

where $V_0(s_0) = -i\omega \mathbf{s}_1(s_0)$ is the Lagrangian velocity of the wall surface element with equilibrium position $s_0$ and instantaneous position $s$. The no-slip boundary condition on the Eulerian fluid velocity $\mathbf{v}(r, t)$ is imposed at the instantaneous surface position $s(t)$,

$$\mathbf{v}(s_0 + s_1 e^{-i\omega t}, t) = V_0(s_0) e^{-i\omega t}, \quad \text{no-slip condition.} \quad (7)$$

Combining Eqs. (5c) and (7) with the Taylor expansion $v_1(s_0 + s_1, t) \approx v_1(s_0) e^{-i\omega t} + \langle (s_1 \cdot \nabla)v_1 \rangle |_{s_0}$, and collecting the terms order by order, gives

$$v_1(s_0) = V_0(s_0), \quad 1\text{st--order condition,} \quad (8a)$$

$$v_2(s_0) = -\langle (s_1 \cdot \nabla)v_1 \rangle |_{s_0}, \quad 2\text{nd--order condition.} \quad (8b)$$

Note that the expansion, or Stokes drift, in Eq. (8b) is valid, if the length scale over which $v_1$ varies is much larger than $|s_1|$. So we require $|s_1| \ll d$ and $|s_2| \ll \delta$.

### D. The limit of weakly curved, thin boundary layers

The crux of our work is the analytical treatment of weakly curved, thin viscous boundary layers. This notion is quantified using the boundary-layer length scale $\delta$ and the compressional length scale $d$,

$$\delta = \sqrt{\frac{2\mu_0}{\omega}}, \quad d = \min \{k_0^{-1}, R\}, \quad (9)$$

where $d$ is the minimum of the wavelength scale $k_0^{-1} = c_0/\omega$ and the length scale $R$ over which the surface curves. We express our subsequent analysis to lowest order in $\epsilon$, defined as the ratio of these length scales,

$$\epsilon = \frac{\delta}{d} \ll 1, \quad (10)$$

where the inequality holds in the limit of weakly curved ($\delta/R \ll 1$), thin boundary layers ($k_0\delta \ll 1$), a condition usually satisfied in microfluidic devices.

### E. Local boundary-layer coordinates

The limit $\epsilon \ll 1$ allows for drastic simplifications of the otherwise complex analytical expressions for curvilinear derivatives of fields inside the boundary layers at distances of order $\delta$ or smaller from the wall. To see this, we introduce the local, right-handed, orthogonal, curvilinear coordinate system with coordinates $\xi$, $\eta$, and $\zeta$. The latter measures distance away from the surface equilibrium position along the surface unit normal vector $\mathbf{e}_n$, while the tangential coordinates $\xi$ and $\eta$ increase in the respective directions of the unit tangent vectors $\mathbf{e}_t$ and $\mathbf{e}_\eta$, but not necessarily measuring arc length, see Fig. 1. To make the scale $R$ of the curved surface explicit, we use the vectorial notation for curvilinear derivatives and introduce the differential-geometric symbols employed in previous boundary-layer analyses in the literature.

For $\epsilon \ll 1$, we may separate any field $A$ inside the boundary layer in the perpendicular coordinate $\zeta$,

$$A(\xi, \eta, \zeta) = A^0(\xi, \eta) a(\zeta), \quad \zeta \equiv \delta \ll d. \quad (12)$$

Here, superscript “0” defines a surface field $A^0(\xi, \eta)$ = $A(\xi, \eta, 0)$, such as the wall velocity $V_0$ and the fluid velocity $v_0$ at the equilibrium position $s_0$ of the wall. Note that a surface field does not have a perpendicular derivative, although it does have a perpendicular component. This coordinate separation results in the following expressions in vectorial notation for the divergence $(A6)$ and advective derivative $(A8)$ involving surface fields:

$$\mathbf{V} \cdot A^0 = \mathbf{V} \cdot A_\parallel^0 + \mathcal{H}_\zeta A_\parallel^0, \quad (13a)$$

$$(A^0 \cdot \mathbf{V}) B^0 = A^0 \cdot (\mathbf{V} B_0^0) e_t + A_\xi^0 B^0_{\zeta} T_{ik} e_i, \quad (13b)$$

$$A_\parallel = A_\xi \mathbf{e}_\xi + A_\eta \mathbf{e}_\eta, \quad (13c)$$

$$\mathbf{V} \cdot \overline{\mathbf{v}}^\parallel = e_\zeta \overline{\partial}_\zeta + e_\eta \overline{\partial}_\eta, \quad (13d)$$

where subscript “$\parallel$” denotes tangential components. See Appendix A for supplemental details.

Importantly, for fluid fields, we distinguish between bulk fields $A^d$ that extend into the bulk with spatial variation on the compressional length scale $d$ and that are typically found by numerical simulation, and boundary-layer fields $A^\delta$ that decays to zero away from the wall at the boundary-layer length scale $\delta$, as sketched in Fig. 1.

$$A^\delta = A^{0\delta}(\xi, \eta) a^\delta(\zeta), \quad \text{with } a^\delta(\zeta) \rightarrow 0 \text{ for } \frac{\zeta}{\delta} \rightarrow \infty. \quad (14)$$

This specific property makes it possible to obtain analytical solutions for the boundary-layer fields $A^\delta$, because the
simplify to derivatives of scalar and vector boundary-layer fields thus

\[ \nabla^2 g^\delta \approx \partial^2 g^\delta, \]

\[ \nabla^2 A^\delta \approx A^\delta \partial^2 A^\delta, \]

\[ \mathbf{V} \cdot A^\delta \approx \mathbf{V} \cdot A^\delta + \mathbf{\partial} A^\delta. \]

With Eqs. (13) and (15), we have established to leading order in \( \epsilon \) the expressions in vectorial form for the curvilinear derivatives in the boundary layer necessary for the subsequent analytical treatment of the boundary layer. In summary, the length-scale conditions for our theory to be valid, in particular Eqs. (8) and (15), are

\[ \delta \ll d, \quad |s_1|| \ll d, \quad |s_1^\perp| \ll \delta. \]

III. FIRST-ORDER TIME-HARMONIC FIELDS

Returning to the perturbation expansion (5), we write the first-order part of the governing equations (2),

\[ p_1 = \epsilon_0^2 p_1, \]

\[ -i\omega \partial_0 p_1 = -\mathbf{V} \cdot \mathbf{v}_1, \]

\[ -i\omega \rho_0 \mathbf{v}_1 = -\mathbf{V}[p_1 - \beta \eta_0 \mathbf{V} \cdot \mathbf{v}_1] + \eta_0 \nabla^2 \mathbf{v}_1, \]

we make a standard Helmholtz decomposition of the velocity field \( \mathbf{v}_1 \),

\[ \mathbf{v}_1 = \mathbf{v}_1^d + \mathbf{v}_1^\delta, \text{ where } \mathbf{V} \times \mathbf{v}_1^d = \mathbf{0} \text{ and } \mathbf{V} \cdot \mathbf{v}_1^d = 0, \]

and insert it into Eq. (17). We separate the equations in solenoidal and irrotational parts and find

\[ i\omega \partial_0 p_1 = \mathbf{V} \cdot \mathbf{v}_1^d, \]

\[ -i\omega \rho_0 \mathbf{v}_1^d = \mathbf{V} \cdot \mathbf{\sigma}_1^d = -(1 - i\Gamma) \mathbf{V} p_1, \]

\[ -i\omega \rho_0 \mathbf{v}_1^\delta = \mathbf{V} \cdot \mathbf{\sigma}_1^\delta = \eta_0 \nabla^2 \mathbf{v}_1^\delta. \]

From this, we derive Helmholtz equations for the bulk fields \( p_1 \) and \( \mathbf{v}_1^d \) as well as for the boundary-layer field \( \mathbf{v}_1^\delta \),

\[ \nabla^2 p_1 + k_v^2 p_1 = 0, \text{ where } k_v = \sqrt{1 + \frac{\Gamma}{2}} k_0, \]

\[ \nabla^2 \mathbf{v}_1^d + k_v^2 \mathbf{v}_1^d = \mathbf{0}, \]

\[ \nabla^2 \mathbf{v}_1^\delta + k_v^2 \mathbf{v}_1^\delta = \mathbf{0}, \text{ where } k_v = \frac{1 + i}{\delta}. \]

Here, we have introduced the compressional wavenumber \( k_v \) in terms of \( \Gamma \) defined in Eq. (4) and \( k_0 = \omega/c_0 \), and the shear wave number \( k_v \) in terms of \( \delta \). Note that \( \Gamma \) is of second order in \( \epsilon \),

\[ \Gamma = \frac{1 + \beta}{2} (k_0 \delta)^2 \sim \epsilon^2 \ll 1. \]

From Eq. (19b) follows that the long-range velocity \( \mathbf{v}_1^d \) is a potential flow proportional to \( \mathbf{V} p_1 \), and as such it is the acoustic velocity of pressure acoustics. The short-range velocity \( \mathbf{v}_1^\delta \) is confined to the thin boundary layer of width \( \delta \) close to the surface, and therefore it is typically not observed in experiments and is ignored in classical pressure acoustics. In the following we derive an analytic solution for the boundary-layer field \( \mathbf{v}_1^\delta \), which is used to determine a boundary condition for \( p_1 \). In this way, the viscous effects from the boundary layer are taken into account in computations of the long-range pressure-acoustic fields \( p_1 \) and \( \mathbf{v}_1^d \).

A. Analytical form of the first-order boundary-layer field

By using Eq. (15b), the analytical solution \( \mathbf{v}_1^d \) to Eq. (20c) is found to be

\[ \mathbf{v}_1^d = \mathbf{v}_1^{d0}(\xi, \eta) e^{ik_\perp \xi} + \mathcal{O}(\epsilon), \]

which describes a shear wave heavily damped over a single wave length, as it travels away from the surface with speed \( c_{sw} = \omega_0 \delta \ll c_0 \). To satisfy the boundary condition (8a), we impose the following condition for \( \mathbf{v}_1^{d0} \) at the equilibrium position \( s_0 \) of the wall,

\[ \mathbf{v}_1^{d0} = \mathbf{V}_1^{d0} - \mathbf{v}_1^{d0}, \text{ first-order no–slip condition}. \]

B. Boundary condition for the first-order pressure field

We now derive a boundary condition for the first-order pressure field \( p_1 \), which takes the viscous boundary layer effects into account without explicit reference to \( \mathbf{v}_1 \). First, it is important to note that the incompressibility condition \( \mathbf{V} \cdot \mathbf{v}_1 = 0 \) used on Eq. (22a) leads to a small perpendicular short-range velocity at \( s_0 \),

\[ \mathbf{v}_1^{d0} = \frac{i}{k_s} \mathbf{V} \cdot \mathbf{v}_1^{d0} = \frac{i}{k_s} \mathbf{V} \cdot \mathbf{V}_1^{d0} - \frac{i}{k_s} \mathbf{V} \cdot \mathbf{v}_1^{d0}. \]

Because \( k_v^{-1} \ll \delta \) and \( \mathbf{V} \cdot \mathbf{v}_1^{d0} \ll a^{-1} \), we find that \( |\mathbf{v}_1^{d0}| \ll |\mathbf{v}_1^d| \), and we repeatedly exploit this relation to neglect terms with \( \mathbf{v}_1^{d0} \) in the following analyses to lowest order in \( \epsilon \).

Using the no-slip condition (22b), the boundary condition on the long-range velocity becomes

\[ \mathbf{v}_1^{d0} = \frac{\mathbf{V}_1^{d0} - \mathbf{v}_1^{d0}}{k_s} \]

\[ = \left( \mathbf{V}_1^{d0} - \frac{i}{k_s} \mathbf{V} \cdot \mathbf{v}_1^{d0} \right) + \frac{i}{k_s} \mathbf{V} \cdot \mathbf{v}_1^{d0} \]

\[ \approx \left( \mathbf{V}_1^{d0} - \frac{i}{k_s} \mathbf{V} \cdot \mathbf{v}_1^{d0} \right) + \frac{i}{k_s} \mathbf{V} \cdot \mathbf{v}_1^{d0}, \]
where the last step is written for later convenience using \((i/k_s) \mathbf{V} \cdot (\mathbf{v}^0 - \mathbf{v}^0) = (i/k_s) (\mathbf{V} \cdot \mathbf{v}^0 - \mathbf{v}^0) = (i/k_s) \mathbf{V} \cdot \mathbf{v}^0\) from Eqs. (13a) and (22b) and using that \(\mathbf{v}^0 = \varepsilon \mathbf{v}_1\). This boundary condition involves the usual expression \(\mathbf{v}^0\) used in classical pressure acoustics plus an \(O(\varepsilon)\)-correction term proportional to \(k_s^{-1}\), due to the parallel divergence of fluid velocity inside the boundary layer that forces a fluid flow perpendicular to the surface to the incompressibility of the short-range velocity \(\mathbf{v}_1\). Note that this correction term is generated partly by the external fluid motion \(-(i/k_s) \mathbf{V} \cdot \mathbf{v}^0\) and partly by the fluid motion itself \((i/k_s) \mathbf{V} \cdot \mathbf{v}^0\). Hence, the wall can affect the long-range fields either by a perpendicular component \(\mathbf{v}^0\), or by a parallel divergence \(\mathbf{V} \cdot \mathbf{v}^0\). The correction term \((i/k_s) \mathbf{V} \cdot \mathbf{v}^0\) due to the fluid motion itself gives the boundary-layer damping of the acoustic energy, see Sec. IV.

Finally, we write Eq. (24b) in the terms of the pressure \(p_1\) using \(\mathbf{V} \cdot \mathbf{v}^0 = \mathbf{V} \cdot \mathbf{v}_1 - \partial_t \mathbf{v}^0\), and Eq. (19),

\[
\partial_t p_1 = \frac{i \omega \rho_0}{1 - i \Gamma} \left( V_{1c}^0 - \frac{i}{k_s} \mathbf{V} \cdot \mathbf{v}^0 \right) - \frac{i}{k_s} \left( k_s^2 p_1 + \partial_t^2 p_1 \right),
\]

boundary condition at \(s_0\).

(25)

C. Boundary condition for the first-order stress

The boundary condition for the first-order stress \(\sigma_1 \cdot \mathbf{e}_z\) on the surrounding wall is found using Eqs. (2c) and (2d). In the viscous stress \(\tau_1\), the divergence terms are neglected, because (19a) leads to \(\eta_0 \mathbf{V} \cdot \mathbf{v}^0 \approx \eta_0 \partial_t \mathbf{v}^0\), and we obtain \(\sigma_1 \cdot \mathbf{e}_z = -p_1 \mathbf{e}_z + \eta_0 \partial_t \mathbf{v}^0\) at \(s_0\). Here, we insert \(\partial_t \mathbf{v}^0 = \partial_t \mathbf{v}_1\) from Eq. (22a), and use Eqs. (19b) and (22b) to express \(\sigma_1 \cdot \mathbf{e}_z\) in terms of the long-range pressure \(p_1\) and wall velocity \(\mathbf{V}^0\) to lowest order in \(\Gamma \sim (k_0 \delta)^2\),

\[
\sigma_1 \cdot \mathbf{e}_z = -p_1 \mathbf{e}_z + i k_s \eta_0 \left( \mathbf{V}^0 + \frac{i}{\omega \rho_0} \mathbf{V} \cdot \mathbf{v}^0 \right),
\]

boundary condition at \(s_0\).

(26)

This is the usual pressure condition plus a correction term of order \(\varepsilon\) due to the viscous shear stress \(\eta_0 \partial_t \mathbf{v}^0\) from the boundary layer.

Equations (20), (24), (25), and (26) constitute our main theoretical result for the first-order acoustic fields. Remarkably, explicit reference to the curvilinear quantities are absent in these equations, only the notion of perpendicular and tangential directions and components are important. In the numerical implementation of them in Sec. VII, we use Cartesian coordinates.

IV. ACOUSTIC POWER LOSS

From the pressure \(p_1\), we derive an expression for the acoustic power loss solely in terms of long-range fields. We introduce the energy density \(E_{ac}^d\) and the energy-flux density \(S_{ac}^d\) of the long-range acoustic fields,

\[
E_{ac}^d(\mathbf{r}, t) = \frac{\mathbf{Re}(p_1 e^{-i\omega t})}{2} + \frac{\rho_0}{2} \mathbf{Re}(\mathbf{v}^d e^{-i\omega t})^2,
\]

with the time averages

\[
\langle E_{ac}^d \rangle = \frac{1}{4} \kappa_0 |p_1|^2 + \frac{\rho_0}{2} |\mathbf{v}^d|^2,
\]

\[
\langle S_{ac}^d \rangle = \langle p_1 \mathbf{v}^d \rangle = \frac{\gamma_0^2}{2} (p_1 \mathbf{v}^d)^2.
\]

In terms of real-valued physical quantities, Eqs. (19a) and (19b) become \(\kappa_0 \partial_t \mathbf{Re}(p_1 e^{-i\omega t}) = -\mathbf{V} \cdot \mathbf{Re}(\mathbf{v}^d e^{-i\omega t}) - \mathbf{p}_0 \mathbf{Re}((\mathbf{v}^d e^{-i\omega t}) = -\mathbf{V} \cdot \mathbf{Re}(1 + i\Gamma) p_1 e^{-i\omega t}).\) Taking the scalar product of \(\mathbf{Re}(\mathbf{v}^d e^{-i\omega t})\) with the latter leads to expressions for the time derivative \(\partial_t E_{ac}^d\) and its time-averaged value \(\langle \partial_t E_{ac}^d \rangle\), which is zero due to the harmonic time dependence,

\[
\partial_t E_{ac}^d = -\mathbf{V} \cdot S_{ac}^d - \Gamma \rho_0 \mathbf{v}^d e^{-i\omega t},
\]

\[
-\mathbf{V} \cdot S_{ac}^d = \frac{1}{2} \Gamma \rho_0 |\mathbf{v}^d|^2/\omega |t|^2.
\]

The latter expression describes the local balance between the convergence of energy-flux density \(S_{ac}^d\) and the rate of change of acoustic energy due to the combined effect of viscous dissipation and viscous energy flux. See Appendix B for a more detailed discussion of this point. Integrating Eq. (29b) over the entire fluid domain \(\Omega\), and using Gauss’s theorem with the \(\zeta\)-direction pointing into \(\Omega\), leads to the global balance of energy rates,

\[
\int_{\partial \Omega} \langle p_1 \mathbf{v}^d \rangle \mathbf{d}A = \frac{1}{2} \Gamma \rho_0 |\mathbf{v}^d|^2 dV.
\]

(30)

This general result reduces to that of classical pressure acoustics only in the special case where \(\mathbf{v}^0 = \mathbf{V}^0\). As seen from Eq. (24c), \(\mathbf{v}^0\) is generated partly externally by the wall motion, and partly internally by the fluid motion. Inserting Eq. (24c) into Eq. (30), and separating wall-velocity terms from fluid-velocity terms gives

\[
\frac{1}{\Omega} \int_{\partial \Omega} \langle p_1 \mathbf{v}^d \rangle \mathbf{d}A = \frac{1}{2} \Gamma \rho_0 |\mathbf{v}^d|^2 dV - \frac{1}{\Omega} \int_{\partial \Omega} \left( \frac{1}{k_s} \mathbf{V} \cdot \mathbf{v}^0 \right) \mathbf{d}A.
\]

(31)

Here, the left-hand side represents the acoustic power gain due to the wall motion, while the right-hand side represents the acoustic power loss \(\langle P_{loss}^d \rangle\) due to the fluid motion. Integrating the last term by parts and using that \(\int_{\partial \Omega} \mathbf{V} \cdot \mathbf{v}^0 \mathbf{d}A = 0\) for any closed surface, we can by Eq. (19b) rewrite \(\langle P_{loss}^d \rangle\) to lowest order in \(\Gamma\) as

\[
\frac{1}{\Omega} \langle P_{loss}^d \rangle = \int_{\Omega} \frac{\Gamma \rho_0 |\mathbf{v}^d|^2 dV + \int_{\partial \Omega} \frac{\delta}{4} \frac{\rho_0 |\mathbf{v}^d|^2}{dA},
\]

which is always positive. The quality factor \(Q\) of an acoustic cavity resonator can be calculated from the long-range fields \(E_{ac}^d\) in Eq. (28a) and \(P_{loss}^d\) in Eq. (32) as
\[ Q = \frac{1}{\omega} \langle \langle E_{\text{sc}} \rangle \rangle \text{d}V \]

We emphasize that in general, \( \langle P_{\text{loss}} \rangle \) is not identical to the viscous heat generation \( \langle P_{\text{loss}}^{\text{visc}} \rangle = \int \langle \nabla v_1 : \tau_1 \rangle \text{d}V \), although as discussed in Appendix B, these might be approximately equal in many common situations.38

\section*{V. SECOND-ORDER STREAMING FIELDS}

As specified in Sec. II B, we only consider the time-averaged streaming and not time-dependent streaming as done by Müller and Bruus.36 For notational simplicity, we therefore drop the angular bracket \( \langle \cdot \rangle \) from the time-averaged velocity \( v_2 \), pressure \( p_2 \), and stress \( \sigma_2 \). The streaming \( v_2 \) is governed by the time-averaged part of Eq. (2) to second order in \( Ma = (1/c_0)|v_1| \), together with the boundary condition (8b),

\[ 0 = \nabla \cdot (\rho_0 v_2 + \langle \rho_1 v_1 \rangle), \quad \text{for} \ r \in \Omega, \quad (34a) \]

\[ 0 = \nabla \cdot \sigma_2 - \rho_0 \nabla \cdot (v_1 v_1), \quad \text{for} \ r \in \Omega, \quad (34b) \]

\[ 0 = v_2 + \langle (s_1 \cdot \nabla) v_1 \rangle, \quad \text{at} \ s_0. \quad (34c) \]

For the given first-order fields \( \rho_1 \) and \( v_1 \), this is a linear Stokes flow problem for \( v_2 \) and \( \sigma_2 \). We decompose the problem into one part driven by the long-range source terms, \( \nabla \cdot (\rho_1 v_1^d) \) in Eq. (34a) and \( \rho_0 \nabla \cdot (v_1^d v_1^d) \) in Eq. (34b), and another part driven by the short-range source terms \( \nabla \cdot (\rho_1 v_1^s) \) and \( \rho_0 \nabla \cdot (v_1^s v_1^s + v_1^d v_1^d + v_1^d v_1^s) \). The corresponding responses are long-range bulk fields \( "d" \) and short-range boundary-layer fields \( "s" \),

\[ v_2 = v_2^d + v_2^s, \quad (35a) \]

\[ p_2 = p_2^d + p_2^s, \quad (35b) \]

\[ \sigma_2 = \sigma_2^d + \sigma_2^s, \quad (35c) \]

\[ v_2^{s0} = -v_2^{d0} - \langle (s_1 \cdot \nabla) v_1 \rangle, \quad \text{at} \ s_0. \quad (35d) \]

Given the boundary conditions Eqs. (35d) and (36d), this length-scale-based decomposition of the linear Stokes problem is unique, see Eqs. (36) and (48), but in contrast to the first-order decomposition (18), it is not a Helmholtz decomposition. Nevertheless, the computational strategy remains the same: we find analytical solutions to the short-range \( \delta \)-fields, and from this we derive boundary conditions for the long-range \( d \)-fields.

Note that our method to calculate the steady second-order fields differs from the standard method of matching “inner” boundary-layer solutions with “outer” bulk solutions.33–35 Our short- and long-range fields co-exist in the boundary layer, but are related by imposing boundary conditions at \( s_0 \).

\section*{A. Short-range boundary-layer streaming}

The short-range part of Eq. (34) consists of all terms containing at least one short-range \( \delta \)-field,

\[ 0 = \nabla \cdot (\rho_0 v_2^d + \langle \rho_1 v_1^d \rangle), \quad (36a) \]

\[ 0 = -\rho_0 \nabla \cdot (v_1^d v_1^d + v_1^d v_1^s + v_1^s v_1^s) + \nabla \cdot \sigma_2^d, \quad (36b) \]

\[ \nabla \cdot \sigma_2^s = \nabla (-p_2^s + \beta \eta_0 \nabla \cdot v_2^s + \eta_0 \nabla^2 v_2^s), \quad (36c) \]

where \( v_2^s \to 0 \) as \( \zeta \to \infty \). \quad (36d)

Notably, condition (36d) leads to a nonzero short-range streaming velocity \( v_2^{s0} \) at the wall, which, due to the full velocity boundary condition (34c), in turn implies a slip condition \( v_2^{s0} \) (35d) on the long-range streaming velocity.

First, we investigate the scaling of \( p_2^s \) by taking the divergence of Eq. (36b) and using Eqs. (36a) and (36c) together with \( \nabla \cdot v_1^d = 0 \) and Eq. (19),

\[ \nabla^2 p_2^s = -\nu_0 (1 + \beta) \nabla^2 (v_1^d \cdot \nabla p_1) \]

\[ -\rho_0 \nabla \cdot \left( \nabla \cdot \langle v_1^d v_1^d + v_1^d v_1^s + v_1^s v_1^s \rangle \right) \]

\[ = -\rho_0 \nabla \cdot \langle 2 v_1^d v_1^d + v_1^d v_1^s + v_1^s v_1^d \rangle \]

\[ -\rho_0 \langle \nabla^2 (2 v_1^d v_1^d + v_1^d v_1^s + v_1^s v_1^d) \rangle. \quad (37b) \]

Recalling from Eq. (23) that \( |v_1^{d0}| \sim \delta d^{-1} v_1 \), we find\( |\rho_0 \langle \nabla^2 v_1^d \rangle : (v_1^d v_1^d)^\dagger \sim (\delta d)^{-1} \rho_0 v_1^2 \) which is the largest possible scaling of the right-hand side. Since by definition \( p_2^s \) is a boundary-layer field, we have \( |\nabla^2 p_2^s| \sim \delta^{-2} p_2^s \), and the scaling of \( p_2^s \) becomes

\[ |p_2^s| \sim \sim \rho_0 v_1^2. \quad (38) \]

Thus, \( \nabla p_2^s \) can be neglected in the parallel component of Eq. (36b), but not necessarily in the perpendicular one. Similarly, in Eq. (36c) we have \( \nabla (\beta \eta_0 \nabla \cdot v_2^s) \]

\[ = -\beta \eta_0 \nabla \cdot \langle v_1^d \rangle \rho_1 \]

which scales as \( \beta \eta_0 \delta^{-2} (v_1^d/c_0) \) and thus much smaller than \( \eta_0 \nabla^2 v_2^s \sim \eta_0 \delta^{-2} (v_1^d/c_0) \).

Henceforth, using the approximation (15b) for the boundary-layer field \( v_2^s \) in Eq. (36b), we obtain the parallel equation to lowest order in \( \epsilon \),

\[ \nu_0 \partial^2_{\zeta} v_2^{s||} = \left| \nabla \cdot (v_1^d v_1^d + v_1^d v_1^s + v_1^s v_1^d) \right|. \quad (39a) \]

Combining this with Eq. (36a), and using Eqs. (15c) and (18), leads to an equation for the perpendicular component \( v_2^{s0} \) of the short-range streaming velocity,

\[ \partial^2_{\zeta} v_2^{s0} = -\nabla_{||} \cdot v_2^{s||} - \frac{1}{\rho_0} \langle v_1^d \cdot \nabla p_1 \rangle. \quad (39b) \]

To determine the analytical solution for \( v_2^{s0} \) in Eq. (39a), we Taylor-expand \( v_1^d \) to first order in \( \zeta \) in the boundary layer, and we use the solution (22a) for \( v_1^d \),

\[ v_1^d = v_1^{d0} + (\partial_\zeta v_1^{d0})^0 \zeta, \quad \text{for} \ \zeta \ll d, \quad (40a) \]

\[ v_1^{d0} = v_1^{d0} q(\zeta), \quad \text{with} \ q(\zeta) = \epsilon^{2k_1} \zeta. \quad (40b) \]

With these expressions, Eq. (39a) becomes
\[ v_{ij} = \frac{1}{2\nu_0} \text{Re} \left\{ \left( \frac{J^{(2)}_{q_1}}{q_{ij}} \right) \nabla \cdot (v_{ij} v_{i0}^0) \right\} + \frac{1}{2\nu_0} \text{Re} \left\{ \left( \frac{J^{(1)}_{q_1}}{q_{ij}} \right) \nabla \cdot (v_{ij} v_{i1}^0) \right\} + \frac{1}{2\nu_0} \text{Re} \left\{ \left( \frac{J^{(3)}_{q_1}}{q_{ij}} \right) \nabla \cdot (v_{ij} v_{i0}^0) \right\} + \frac{k_0}{2\rho_0} \text{Re} \left\{ \left( \frac{J^{(3)}_{q_1}}{q_{ij}} \right) \nabla \cdot (v_{ij} v_{i0}^0) \right\}. \tag{44} \]

We have neglected the term \((1/2\nu_0)\text{Re} \left\{ \left( \frac{J^{(1)}_{q_1}}{q_{ij}} \right) \nabla \cdot (v_{ij} v_{i0}^0) \right\} \) as \(v_{ij} \sim c_0 v_{i0}^0\), due to Eq. (23), and the two terms proportional to \(J_{q_1}^{(1)}\) and \(J_{q_1}^{(2)}\), as these are \(\sim \delta^3\). Remarkably, the term \(J_{q_1}^{(1)} v_{ij} v_{i0}^0 v_{i1}^0\) may scale with an extra factor \(e^{-1}\) compared to all other terms, and thus may dominate the boundary-layer velocity. However, in the computation of the long-range slip velocity \(v_{ij}^{\delta 0}\) in Sec. V.B, its contribution is canceled by the Stokes drift \(s_1 \cdot \nabla v_{i1}\), as also noted in Ref. 35. Using \(v_{ij}^{\delta 0} = V_{ij} - v_{ij}^0\), the property \(I^{(n)}_{ab} = I_{11}^{(n)}\), and rearranging terms, we arrive at

\[ v_{ij}^{\delta 0} = \frac{1}{2\nu_0} \text{Re} \left\{ \left( \frac{J^{(2)}_{q_1}}{q_{ij}} \right) \nabla \cdot (v_{ij} v_{i0}^0) \right\} + \frac{1}{2\nu_0} \text{Re} \left\{ \left( \frac{J^{(1)}_{q_1}}{q_{ij}} \right) \nabla \cdot (v_{ij} v_{i1}^0) \right\} + \frac{1}{2\nu_0} \text{Re} \left\{ \left( \frac{J^{(3)}_{q_1}}{q_{ij}} \right) \nabla \cdot (v_{ij} v_{i0}^0) \right\} + \frac{k_0}{2\rho_0} \text{Re} \left\{ \left( \frac{J^{(3)}_{q_1}}{q_{ij}} \right) \nabla \cdot (v_{ij} v_{i0}^0) \right\}. \tag{45} \]

The perpendicular short-range velocity component \(v_{ij}^{\delta 0}\) is found by integrating Eq. (39b) with respect to \(\zeta\). The integration of the \(V_{ij} \cdot v_{ij}^{\delta 0}\)-term is carried out by simply increasing the superscript of the \(I^{(n)}_{ab}\)-integrals in Eq. (45) from \((n)\) to \((n+1)\), while the integration of the \(V_{ij} v_{ij}^{\delta 0}\)-term is carried out by using Eq. (19b) to substitute \((1/p_0)\nabla p_1\) by \(io_0 c_0^2 v_i\) and introducing the suitable \(I^{(n)}_{ab}\)-integral for the factor \(q(\zeta)\), namely, \(I^{(1)}_{q_1} = -iI^{(1)}_{q_1}\);
and
\[
v_{l=0}^2 = -\frac{\delta}{2\omega} \text{Re} \left[ \mathbf{v}_l \cdot \left( -\frac{5}{4} \mathbf{V} \cdot (v_{l=0}^0 v_{l=0}^0) + \frac{1-i}{2\omega} \mathbf{V} \cdot (v_{l=0}^0 v_{l=0}^0) + \frac{1+i}{\delta} \mathbf{V} \cdot (v_{l=0}^0 v_{l=0}^0) + \frac{i}{\delta} \mathbf{V} \cdot (v_{l=0}^0 v_{l=0}^0) \right)_{||} - k_0^2 (1-i) v_{l=1}^0 v_{l=0}^0 \right] \tag{47b}
\]
\[
= -\frac{1}{2\omega} \text{Re} \left[ \mathbf{v}_l \cdot \left( i v_{l=1}^0 \mathbf{V} v_{l=0}^0 \right) \right] + O(\epsilon). \tag{47c}
\]

B. Long-range bulk streaming

The long-range part of Eq. (34) is
\[
0 = \mathbf{V} \cdot \left[ \rho_0 v_{l=0}^d + (\rho_1 v_{l=1}^d) \right], \tag{48a}
\]
\[
0 = -\rho_0 \mathbf{V} \cdot \langle v_{l=1}^d v_{l=1}^d \rangle + \mathbf{V} \cdot \sigma_{l=1}^d, \tag{48b}
\]
\[
\mathbf{V} \cdot \sigma_{l=1}^d = -\mathbf{V} \rho_0 \beta \eta_0 \mathbf{V} \cdot v_{l=1}^2 + \eta_0 \nabla^2 v_{l=1}^2, \tag{48c}
\]
\[
v_{l=0}^2 = -v_{l=0}^2 - \langle \mathbf{s}_{l=1} \cdot \mathbf{V} \rangle v_{l=1}, \tag{48d}
\]

In contrast to the limiting-velocity matching at the outer edge of the boundary layer done by Nyborg,\textsuperscript{33} we define the boundary condition (48d) on the long-range streaming \(v_{l=1}^2\) at the equilibrium position \(s_0\).

To simplify Eq. (48), we investigate the products of first-order fields. In Eq. (48a), we use Eq. (29b) and find
\[
\mathbf{V} \cdot v_{l=1}^d = \frac{\langle \rho_1 v_{l=1}^d \rangle}{\rho_0} = -\mathbf{V} \langle \mathbf{S}_{l=1}^d \rangle \frac{\Gamma_k k_0 |v_{l=1}^0|^2}{2\omega} \tag{49}
\]

Since each term in \(\mathbf{V} \cdot v_{l=1}^d\) scales as \((k_0/\epsilon_0)|v_{l=1}^0|^2 \gg (\Gamma/2)(k_0/\epsilon_0)|v_{l=1}^0|^2\), we conclude that \(\mathbf{V} \cdot v_{l=1}^2 \approx 0\) is a good approximation, corresponding to ignoring the small viscous dissipation in the energy balance expressed by Eq. (29b). In Eq. (48b), the divergence of momentum flux can be rewritten using Eq. (19b),
\[
\rho_0 \mathbf{V} \cdot \langle v_{l=1}^d v_{l=1}^d \rangle = -\mathbf{V} \langle \mathbf{L}_{l=1}^d \rangle - \frac{\Gamma \omega}{\epsilon_0} \langle \mathbf{S}_{l=1}^d \rangle, \tag{50}
\]

where we introduced the long-range time-averaged acoustic Lagrangian density,
\[
\langle \mathbf{L}_{l=1}^d \rangle = \frac{1}{4} \kappa_0 |p_{l=1}|^2 - \frac{1}{4\rho_0} |v_{l=1}^0|^2. \tag{51}
\]

Note that \(|\mathbf{V} \langle \mathbf{L}_{l=1}^d \rangle| \sim \omega p_{l=1}^2/\rho_0 c_1^3\), whereas \(|(\Gamma \omega/\epsilon_0^2) \langle \mathbf{S}_{l=1}^d \rangle| \sim \Gamma \omega p_{l=1}^2/(\rho_0 c_1^3)\), so the first term in Eq. (50) is much larger than the second term. However, as also noted by Riaud et al.,\textsuperscript{39} since the first term is a gradient, it cannot drive any rotating streaming. In practice, it is therefore advantageous to work with the excess pressure \(p_{l=1}^2 - \langle \mathbf{L}_{l=1}^d \rangle\). Finally, in Eq. (48c), we again use \(\mathbf{V} \cdot v_{l=1}^2 \approx 0\). With these considerations, Eqs. (48) become those of an incompressible Stokes flow driven by the body force \((\Gamma \omega/\epsilon_0^2) \langle \mathbf{S}_{l=1}^d \rangle\) and the slip velocity \(v_{l=1}^0\) at the boundary,
\[
0 = \mathbf{V} \cdot v_{l=1}^0, \tag{52a}
\]
\[
0 = -\mathbf{V} \left[ \rho_1^0 - \langle \mathbf{L}_{l=1}^d \rangle \right] + \eta_0 \nabla^2 v_{l=1}^0 + \frac{\Gamma \omega}{\epsilon_0} \langle \mathbf{S}_{l=1}^d \rangle, \tag{52b}
\]
\[
v_{l=1}^0 = -v_{l=0}^0 - \langle \mathbf{s}_{l=1} \cdot \mathbf{V} \rangle v_{l=1}, \tag{52c}
\]

These equations describe acoustic streaming in general. The classical Eckart streaming\textsuperscript{40} originates from the body force \((\Gamma \omega/\epsilon_0^2) \langle \mathbf{S}_{l=1}^d \rangle\), while the classical Rayleigh streaming\textsuperscript{32} is due to the boundary condition (52c).

The Stokes drift \(\langle \mathbf{s}_{l=1} \cdot \mathbf{V} v_{l=1} \rangle |_{s_0}\), induced by the oscillating wall, is computed from Eqs. (6), (18), and (22a),
\[
\langle \mathbf{s}_{l=1} \cdot \mathbf{V} v_{l=1} \rangle |_{s_0} = -\frac{1}{2\omega} \text{Re} \left[ i V_{l=1}^{00} \cdot \mathbf{V} \left( v_{l=1}^0 + v_{l=0}^0 \right) \right] |_{s_0} \tag{53}
\]

From this, combined with Eqs. (47) and (52c), follows the boundary condition \(v_{l=0}^0\) for the long-range streaming velocity \(v_{l=1}^2\) expressed in terms of the short-range velocity \(v_{l=1}^2\) and the wall velocity \(V_{l=1}^0\). The parallel component is
\[
v_{l=1}^0 = -\frac{1}{2\omega} \text{Re} \left\{ \mathbf{V} \cdot \left[ \frac{i}{2} v_{l=1}^0 v_{l=1}^{00} + i v_{l=1}^{00} v_{l=1}^{00} \right] - i V_{l=1}^{00} \cdot \mathbf{V} \left( v_{l=1}^0 + v_{l=0}^0 \right) \right\}, \tag{54a}
\]

where the large terms proportional to \(\langle (1+i/\delta) V_{l=1}^{00} v_{l=1}^0 \rangle\) canceled out, as also noted by Vanneste and Bühler.\textsuperscript{35} Similarly, the perpendicular component becomes
\[
v_{l=1}^{00} = \frac{\delta}{2\omega} \text{Re} \left[ -k_0^2 (1-i) v_{l=0}^{00} \cdot v_{l=1}^{00} \right] - \mathbf{V} \cdot \langle \mathbf{L}_{l=1}^{00} \rangle - \frac{\Gamma \omega}{\epsilon_0} \langle \mathbf{S}_{l=1}^{00} \rangle, \tag{54b}
\]

\[
= \frac{1}{2\omega} \text{Re} \left[ \mathbf{V} \cdot \left( i v_{l=1}^{00} V_{l=1}^{00} \right) - \frac{1+i}{\delta} V_{l=1}^{00} v_{l=1}^{00} \right] + O(\epsilon). \tag{54c}
\]
Taking the divergences in Eq. (54a) and using Eq. (23), as well as computing Eq. (54c) to lowest order in $\epsilon$, leads to the final expression for the slip velocity at $\zeta = 0$,

$$v_{20} = (A \cdot e_x)e_x + (A \cdot e_y)e_y + (B \cdot e_z)e_z,$$

$$A = -\frac{1}{2\omega} \text{Re} \left\{ v_{10}^{00} \cdot \mathbf{V} \left( \frac{1}{2} v_{10}^{00} - i \mathbf{V}_1^{0} \right) - i \mathbf{V}_1^{0} \cdot \nabla v_{10}^{00} \right\},$$

$$B = \frac{1}{2\omega} \text{Re} \left\{ i v_{10}^{00} \cdot \nabla v_{10}^{00} \right\},$$

(55)

where $A$ and $B$ are associated with the parallel and perpendicular components $v_{10}^{00}$ and $v_{20}^{00}$, respectively, and where to simplify we used $(v_{10}^{00} \cdot \mathbf{V}) v_{10}^{00} = v_{10}^{00} v_{10}^{00}$ and the relations $T_{3i} = 0$, $T_{kij} = 0$, and $T_{ikj} = T_{jki}$ for the curvilinear quantities, see Eq. (A3) in Appendix A.

Equations (52) and (55) constitute our main theoretical result for the second-order acoustic streaming.

VI. SPECIAL CASES

In the following, we study some special cases of our main results (20a) and (25) for the acoustic pressure $p_1$ and Eqs. (52) and (55) for the streaming velocity $v_2^0$, and relate them to previous studies in the literature.

A. Wall oscillations restricted to the perpendicular direction

The case of a weakly curved wall oscillating only in the perpendicular direction was studied by Nyborg and later refined by Lee and Wang. Using our notation, the boundary conditions used in these studies were

$$v_{10}^{0} + v_{10}^{00} = V_{10}^{0} = V_{10}^{0} - i_e,$$

(56)

whereby $\mathbf{V} \cdot V_{10}^{0} = \mathcal{H}_{1} V_{10}^{0}$, so that our boundary condition (25) for $p_1$ to lowest order in $\Gamma$ becomes

$$\partial_\parallel p_1 = i \omega p_0 \left( 1 - \frac{1}{k_y} \mathcal{H}_{1} \right) V_{10}^{0} - \frac{i}{k_y} \left( k_x^2 p_1 + \partial_{k_x}^2 p_1 \right).$$

(57)

Similarly, for the steady streaming $v_{2i}^{0}$, Eq. (56) gives $\mathbf{V} \cdot v_{2i}^{0} \approx -\mathbf{V} \parallel \cdot v_{10}^{0} = -\mathbf{V} \cdot v_{10}^{00} = -\mathbf{V} \cdot v_{10}^{00} - \partial_{k_y} v_{10}^{00} - \partial_{k_y} v_{10}^{00}$ evaluated at $\zeta = 0$. Combining this expression with the derivative rule (13b) and the index notation $\zeta = \eta$ and $\eta = \zeta$, as well as $x, \beta = \zeta, \eta$, the boundary condition (55) gives to lowest order in $\epsilon$ the tangential components

$$v_{20}^{0} = -\frac{1}{4\omega} \text{Re} \left\{ v_{10}^{00} \partial_{x} v_{10}^{00} + i v_{10}^{00} v_{10}^{00} T_{3j} + 2i v_{10}^{00} \partial_{x} v_{10}^{00} + \left( 2i + 1 \right) \mathbf{V} \cdot v_{10}^{00} - \left( 2 - 3i \right) \partial_{x} v_{10}^{00} \right\},$$

(58a)

and the perpendicular component

$$v_{20}^{0} = \frac{1}{2\omega} \text{Re} \left\{ i v_{10}^{00} \partial_{k_x} v_{10}^{00} \right\}.$$

(58b)

When comparing our expressions with the results of Lee and Wang, denoted by a superscript “LW” below, we note the following. Neither the pressure $p_1$ nor the steady perpendicular streaming velocity $v_{2i}^{0}$ were studied by Lee and Wang, so our results Eqs. (57) and (58b) for these fields represent an extension of their work. The slip condition (58a) for the parallel streaming velocity $v_{2i}^{0}$ with $\beta = \zeta, \eta$ is presented in Eqs. (19)LW and (20)LW as the limiting values $u_L$ and $v_L$ for the two parallel components of $v_2^0$ outside the boundary layer. A direct comparison is obtained by (1) identifying our $v_{2i}^{0}$ with the acoustic velocity $(u_{00}, v_{00}, w_{00})$LW, and our $T_{jki}$ with $T_{jki}$LW; (2) taking the complex conjugate of the argument of the real value in Eq. (58a), and (3) noting that $q_k$ and $q_\perp$ defined in Eqs. (3)LW and (4)LW equal the first two terms of Eq. (58a). By inspection we find agreement, except that Lee and Wang are missing the terms $2iV_{10}^{0} \left( \partial_{x} v_{10}^{00} + v_{10}^{00} T_{3j} \right)$, which in our calculation partly arise from the Lagrangian velocity boundary condition (34c), where Lee and Wang have used the no slip condition $v_2 = 0$. For more details see Appendix C1.

B. A flat wall oscillating in any direction

The case of a flat wall oscillating in any direction was studied by Vanneste and Bühler. In this case, we adapt Cartesian coordinates $(\xi, \eta, \zeta) = (x, y, z)$, for which all scale factors $h_i$ are unity, $\partial_i = \partial_i$, and all curvilinear quantities $T_{kij}$ and $\mathcal{H}_k$ are zero. The resulting boundary conditions (25) and (55) for the pressure $p_1$ and for the long-range streaming $v_2^0$, then simplify to

$$\partial_\parallel p_1 = i \omega p_0 V_{10}^{0} - \frac{1}{2} \delta \left( \partial_{x} p_0 \mathbf{V} \parallel \cdot V_{10}^{0} + k_1^2 p_1 + \partial_{k_1}^2 p_1 \right),$$

(59a)

$$v_{20}^{0} = -\frac{1}{4\omega} \text{Re} \left\{ \left( 1 - 2i \right) v_{10}^{00} \partial_{x} v_{10}^{00} - 4i v_{10}^{00} \partial_{x} v_{10}^{00} + \left( 2 + i \right) \partial_{x} v_{10}^{00} + 2i \partial_{x} v_{10}^{00} \right\} \nabla v_{10}^{0} - 2i v_{10}^{00} \partial_{k_1} v_{10}^{00} \right\},$$

(59b)

$$v_{20}^{0} = -\frac{1}{4\omega} \text{Re} \left\{ -2i v_{10}^{00} \partial_{k_1} v_{10}^{00} \right\}. $$

(59c)

The pressure condition (59a) was not studied in Ref. 35, so it represents an extension of the existing theory. On the other hand, Eqs. (59b) and (59c) are in full agreement with Eq. (4.14) in Vanneste and Bühler. To see this, we identify our first-order symbols with those used in Ref. 35 as $v_1^0 = 2\mathbf{V} \parallel$ and $v_2^0 = -2U \xi_t - 2V \eta_t$, and we relate our steady Eulerian second-order long-range velocity $v_2^0$ with their Lagrangian mean flow $\bar{u}$ using the Stokes drift expression (34c) as $v_2^0 + (1/\omega)(iv_1^0 \cdot \mathbf{V} \parallel) = \bar{u}$ at the interface $z = 0$. For more details see Appendix C2.

C. Small surface velocity compared to the bulk velocity

At resonance in acoustic devices with a large resonator quality factor $Q \gg 1$, the wall velocity $V_1^0$ is typically a
factor $Q$ smaller than the bulk fluid velocity $v_i^d$, \cite{25,36} which is written as $V_i^0 \sim Q^{-1} v_i^d \ll v_i^f$. In this case, as well as for rigid walls, we use $V_i^0 = 0$ in Eq.\ (55), so that $v_i^0 = -v_i^0$ and

$$\langle \mathbf{v}_i^0 \cdot \mathbf{v}_i^0 \rangle = \langle \mathbf{v}_i^0 \cdot \mathbf{v}_i^d \rangle \approx \frac{1}{4} \left| \mathbf{v}_i^0 \right|^2. \quad (60)$$

Here, $\mathbf{v}_i^0$ is neglected because $\left| \mathbf{v}_i^0 \right| \approx \left| \mathbf{V}_i^0 \right| \ll \left| \mathbf{v}_i^d \right|$, and we have used that $\mathbf{V} \times \mathbf{v}_i^0 = 0$ from Eq.\ (18). Hence, for devices with rigid walls $V_i^0 = 0$, or resonant devices with $\left| \mathbf{v}_i^0 \right| \gg \left| \mathbf{V}_i^0 \right|$, the slip-velocity $v_i^0$ becomes

$$v_i^0 = -\frac{1}{8\omega} \mathbf{V}_i \left| \mathbf{v}_i^d \right|^2 - \text{Re} \left\{ \left( \frac{2-4i}{4\omega} \mathbf{V}_i \cdot \mathbf{v}_i^0 + \frac{i}{2\omega} \delta v_i^d \right) \mathbf{v}_i^0 \right\}, \quad (61a)$$

$$v_i^0 = 0. \quad (61b)$$

Two important limits are parallel acoustics, where $\left| \partial v_i^d \right| \ll \left| \mathbf{V}_i \cdot \mathbf{v}_i^0 \right|$, and perpendicular acoustics, where $\left| \partial v_i^d \right| \gg \left| \mathbf{V}_i \cdot \mathbf{v}_i^0 \right|$. In the first limit, the pressure is mainly related to the parallel velocity variations, and from Eqs. \(19a\) and \(19b\) we have $\mathbf{V}_i \cdot \mathbf{v}_i^0 = i\omega \kappa_0 p_i$ and $\mathbf{v}_i^0 = -(i/\rho_0 \omega) \mathbf{V}_i p_i$. For parallel acoustics we can therefore write Eq. \(61a\) as,

$$v_i^0 = \frac{1}{8\omega \rho_i} \mathbf{V}_i \left[ 2\kappa_0 \left| p_i \right|^2 - \rho_0 \left| \mathbf{v}_i^0 \right|^2 \right] + \kappa_0 / 2 \langle \mathbf{S}^{\perp}_i \rangle,$$

for parallel acoustics, $\left| \partial v_i^d \right| \ll \left| \mathbf{V}_i \cdot \mathbf{v}_i^0 \right|$. \quad (62a)

The classical period-doubled Rayleigh streaming, \cite{22} which arises from a one-dimensional parallel standing wave, results from the gradient-term in Eq.\ (62a). This is seen by considering a rigid wall in the $x$-$y$ plane with a standing wave above it in the $x$ direction of the form $v_i^d = v_{i0} \cos(k_0 x) \hat{e}_x$, where $v_{i0}$ is a velocity amplitude. Inserting this into Eq. \(62a\) yields Rayleigh’s seminal boundary velocity $v_i^0 = (3/8)(v_{i0}^2 / \omega) \sin(2k_0x) \hat{e}_x$. Another equally simple example of parallel acoustics is the boundary condition generated by a planar travelling wave of the form $v_i^d = v_{i0} e^{i(k_0 x + \phi)} \hat{e}_x$.

Here, only the energy-flux density $\langle \mathbf{S}^{\perp}_i \rangle$ in Eq. \(62a\) contributes to the streaming velocity which becomes the constant value $v_i^0 = (1/4)(v_{i0}^2 / \omega) \hat{e}_x$.

The opposite limit is perpendicular acoustics, where the pressure is mainly related to the perpendicular velocity variations $\partial v_i^d \mathbf{v}_i^0 = i\omega \kappa_0 p_i$. In this limit, Eq. \(61a\) is given by a single term

$$v_i^0 = -\kappa_0 \langle \mathbf{S}^{\perp}_i \rangle,$$

for perpendicular acoustics, $\left| \partial v_i^d \right| \gg \left| \mathbf{V}_i \cdot \mathbf{v}_i^0 \right|$. \quad (62b)

We emphasize that in these two limits, the only mechanism that can induce a streaming slip velocity, which rotates parallel to the surface, is the energy-flux density $\langle \mathbf{S}^{\perp}_i \rangle$. As seen from Eq. \(52b\), this mechanism also governs the force density driving streaming in the bulk. In general, $\langle \mathbf{S}^{\perp}_i \rangle$ can drive rotating streaming, if it has a nonzero curl. This we compute to lowest order in $\Gamma$ using Eq. \(19b\) and $\mathbf{V} \times \mathbf{v}_i^d = 0$, and find it to be proportional to the acoustic angular momentum density,

$$\mathbf{V} \times \langle \mathbf{S}^{\perp}_i \rangle = \omega^2 \langle \mathbf{r}^d_i \times (\rho_0 \mathbf{u}_i^d) \rangle, \quad \mathbf{r}^d_i = \frac{i}{\omega} \mathbf{v}_i^d. \quad (63)$$

**VII. NUMERICAL MODELING IN COMSOL**

In the following we implement our extended acoustic pressure theory, Eqs. \(20a\) and \(25\) for $p_i$, and streaming theory, Eqs. \(52\) and \(55\) for $v_i^d$ and $p_2$, in the commercial finite-element-method (FEM) software COMSOL MULTIPHYSICS. \cite{41} We compare these simulations with a full boundary-layer-resolved model for the acoustics, Eqs. \(17\) and \(8a\) for $v_1$ and $p_1$, and for the streaming, Eqs. \(34\) and \(8b\) for $v_2$ and $p_2$. The full model is based on our previous acoustofluidic modeling of fluids-only systems \cite{28,36,42} and solid-fluid systems. \cite{43}

Remarkably, our extended (effective) acoustic pressure model makes it possible to simulate acoustofluidic systems not accessible to the brute-force method of the full model for three reasons: (1) In the full model, the thin boundary layers need to be resolved with a fine FEM mesh. This is not needed in our effective model. (2) For the first-order acoustics, the full model is based on the vector field $v_1$ and the scalar field $p_1$, whereas our effective model is only based on the scalar field $p_1$. (3) For the second-order streaming, the full equations \(34\) contain large canceling terms, which have been removed in the equations \(52\) used in the effective model. Therefore, also in the bulk, the effective model can be computed on a much coarser FEM mesh than the full model.

In Sec. \(VIII\), we model a fluid domain $\Omega_f$ driven by boundary conditions applied directly on $\partial \Omega_f$, and in Sec. \(IX\), we model a fluid domain $\Omega_f$ embedded in an elastic solid domain $\Omega_s$ driven by boundary conditions applied on the outer part of the solid boundary $\partial \Omega_s$.

In \textsc{comsol}, we specify user-defined equations and boundary conditions in weak form using the PDE mathematics module, and we express all vector fields in Cartesian coordinates $(x, y, z)$. At the boundary $\partial \Omega_f$, the local right-handed orthonormal basis $(\hat{e}_x, \hat{e}_y, \hat{e}_z)$ is implemented using the built-in \textsc{comsol} tangent vectors $t_1$ and $t_2$ as well as the normal vector $n$, all given in Cartesian coordinates. Boundary-layer fields (superscript "0"), such as $V_i^0, v_i^0$, and $p_i$, are defined on the boundary $\partial \Omega_f$ only, and their spatial derivatives are computed using the built-in tangent-plane derivative operator $dtang$. For example, in \textsc{comsol} we call the Cartesian components of $v_i^0$ for $vdX, vdy$, and $vdz$ and compute $\mathbf{V} \cdot v_i^0$ as $dtang(vdX, x) + dtang(vdY, y) + dtang(vdZ, z)$. The models are implemented in \textsc{comsol} using the following two-step procedure. \cite{36}

Step (1), first-order fields: \cite{42,43} For a given frequency $\omega$, the driving first-order boundary conditions for the system are specified: the wall velocity $V_i^0$ on $\partial \Omega_f$ for the fluid-only model, and the outer wall displacement $u_i$ on $\partial \Omega_d$ for the solid-fluid model. Then, the first-order fields are solved; the pressure $p_1$ in $\Omega_f$ using Eqs. \(20a\) and \(25\), and, if included in the model, the solid displacement $u_1$ in the solid domain $\Omega_s$.
In particular, in COMSOL we implement \( \partial_\nu p_1 = (e_z \cdot \nabla)^2 p_1 \) in Eq. (25) as \( \nabla \times \nabla \times \mathbf{p}_1 + 2 \nabla \times \mathbf{p}_1 \times \nabla \times \mathbf{p}_1 + \cdots \).

Step (2), second-order fields,\textsuperscript{36,42} Time averages \( \frac{1}{a} \mathbb{E} \{ \mathbf{f}^g \} \) are implemented using the built-in COMSOL operator \( \mathbb{E} \{ \cdot \} \) as \( 0.5 \star \mathbb{E} \{ f \} \). Moreover, in the boundary condition \( \text{(55)} \), the normal derivative of \( \mathbf{v}^1_{ij} \) in \( \mathbf{A} \) is rewritten as \( \partial_n \mathbf{v}^1_{ij} = \mathbf{V} \cdot \mathbf{v}^1_{ij} - \mathbf{v}^1_{ij} = \mathbf{e}_1 \partial_1 \mathbf{v}^1_{ij} = \mathbf{e}_1 \cdot \nabla \mathbf{v}^1_{ij} \) for computational ease, and the advective derivatives in \( \mathbf{A} \) and \( \mathbf{B} \), such as the term \( \mathbb{E} \{ \mathbf{V} \cdot \mathbf{v}^1_{ij} \} \cdot \mathbf{e}_s \), in \( \mathbf{A} \cdot \mathbf{e}_s \) are computed as \( \mathbb{E} \{ \mathbf{v}^1_{ij} \} \mathbb{E} \{ \mathbf{d} \} \). The 2D rectangular fluid domain \( \Omega_2 \) is the full boundary-layer-resolved simulation developed by the authors in Ref. 42.

### VIII. EXAMPLE I: A RECTANGULAR CAVITY

We apply our theory to a long, straight channel along the \( x \) axis with a rectangular cross section in the vertical \( y \)-\( z \) plane, a system intensively studied in the literature both theoretically\textsuperscript{28,36,42} and experimentally.\textsuperscript{25,45-47} We consider the 2D rectangular fluid domain \( \Omega_2 \) with \( -\frac{1}{2} L \leq y < \frac{1}{2} L \) and \( -\frac{1}{2} H < z < \frac{1}{2} H \), where the top and bottom walls at \( z = \pm \frac{1}{2} H \) are stationary and the vertical side walls at \( y = \pm \frac{1}{2} L \) oscillate with a given velocity \( V_{0y}^0 \) and frequency \( \nu = \omega/2\pi \) close to \( c_{02}/2W \), thus exciting a half-wave resonance in the \( y \)-direction. In the simulations we choose the wall velocity to be \( V_{0y}^0 = d_0 \omega \) with a displacement amplitude \( d_0 \approx 0.1 \text{mm} \). The material parameters used in the model are shown in Table I.

We compare the results from the effective theory with the full boundary-layer-resolved simulation developed by Muller et al.\textsuperscript{28} Moreover, we derive analytical expressions for the acoustic fields, using pressure acoustics and our boundary condition using Eq. (55). We apply our theory to a long, straight channel along the \( x \) axis with a rectangular cross section in the vertical \( y \)-\( z \) plane, a system intensively studied in the literature both theoretically\textsuperscript{28,36,42} and experimentally.\textsuperscript{25,45-47} We consider the 2D rectangular fluid domain \( \Omega_2 \) with \( -\frac{1}{2} L \leq y < \frac{1}{2} L \) and \( -\frac{1}{2} H < z < \frac{1}{2} H \), where the top and bottom walls at \( z = \pm \frac{1}{2} H \) are stationary and the vertical side walls at \( y = \pm \frac{1}{2} L \) oscillate with a given velocity \( V_{0y}^0 \) and frequency \( \nu = \omega/2\pi \) close to \( c_{02}/2W \), thus exciting a half-wave resonance in the \( y \)-direction. In the simulations we choose the wall velocity to be \( V_{0y}^0 = d_0 \omega \) with a displacement amplitude \( d_0 \approx 0.1 \text{mm} \). The material parameters used in the model are shown in Table I.

We compare the results from the effective theory with the full boundary-layer-resolved simulation developed by Muller et al.\textsuperscript{28} Moreover, we derive analytical expressions for the acoustic fields, using pressure acoustics and our boundary condition using Eq. (55).

#### A. First-order pressure

To leading order in \( \epsilon \) and assuming small variations in \( z \), Eqs. (20a) and (25) in the fluid domain \( \Omega_2 \) become

\[
\begin{align*}
\nabla^2 p_1 + k_0^2 p_1 &= 0, \quad r \in \Omega_2, \tag{64a}
\end{align*}
\]

### TABLE I. Material parameters at 25 °C used in the numerical modeling presented in Secs. VIII and IX.

<table>
<thead>
<tr>
<th>Material</th>
<th>Property</th>
<th>Value 1</th>
<th>Value 2</th>
</tr>
</thead>
<tbody>
<tr>
<td>Water (Ref. 42)</td>
<td>Mass density</td>
<td>( \rho_0 )</td>
<td>997.05</td>
</tr>
<tr>
<td></td>
<td>Compressibility</td>
<td>( \kappa_0 )</td>
<td>452.0</td>
</tr>
<tr>
<td></td>
<td>Speed of sound</td>
<td>( c_0 )</td>
<td>1496.7</td>
</tr>
<tr>
<td>Pyrex glass</td>
<td>Dynamic viscosity</td>
<td>( \eta_0 )</td>
<td>0.890</td>
</tr>
<tr>
<td></td>
<td>Bulk viscosity</td>
<td>( \eta_0^b )</td>
<td>2.485</td>
</tr>
<tr>
<td></td>
<td>Density</td>
<td>( \rho_d )</td>
<td>2230.0</td>
</tr>
<tr>
<td></td>
<td>Speed of sound, longitudinal</td>
<td>( c_{00} )</td>
<td>5592.0</td>
</tr>
<tr>
<td></td>
<td>Speed of sound, transverse</td>
<td>( c_{tt} )</td>
<td>3424.0</td>
</tr>
</tbody>
</table>

Note that at resonance, the horizontal velocity component is amplified by a factor \( k_0^2 \) since the presence of the boundary layers introduces a small variation in the \( z \) direction. The half-wave resonance that maximizes the amplitude of \( p_1 \) in Eq. (65) is therefore found at a frequency \( f_{\text{res}} \) slightly lower than \( f_{\text{res}}^0 = c_{02}/2W \),

\[
f_{\text{res}} = \left( 1 - \frac{1}{2} \Gamma_\text{bl} \right) f_{\text{res}}^0, \quad \text{with} \quad \Gamma_\text{bl} = \frac{\delta}{
\]

Here, we introduced the boundary-layer damping coefficient \( \Gamma_\text{bl} \) that shifts \( f_{\text{res}} \) away from \( f_{\text{res}}^0 \). This resonance shift arises from the extended boundary condition (25) and is thus beyond classical pressure acoustics.

Using \( f = f_{\text{res}} \) in Eq. (65) and expanding to leading order in \( \Gamma_\text{bl} \), gives the resonance pressure and velocity,

\[
\begin{align*}
\frac{p_{\text{res}}}{\rho_0 c_0} &= -\frac{4 V_{0y}^0}{\pi \Gamma_{\text{bl}}} \left\{ \sin(\tilde{y}) + \frac{\Gamma_{\text{bl}}}{2} i \tilde{y} \cos(\tilde{y}) \right\} Z_{\text{res}}(\tilde{z}), \tag{69a}
\end{align*}
\]

\[
\begin{align*}
v_{y,\text{res}} &= \frac{4 i V_{0y}^0}{\pi \Gamma_{\text{bl}}} \left\{ \cos(\tilde{y}) - \frac{\Gamma_{\text{bl}}}{2} i \tilde{y} \sin(\tilde{y}) \right\} Z_{\text{res}}(\tilde{z}), \tag{69b}
\end{align*}
\]

\[
\begin{align*}
v_{z,\text{res}} &= \frac{4 V_{0y}^0}{\pi} \left( 1 + i \tilde{y} \sin(\tilde{y}) \right) \tilde{z}, \tag{69c}
\end{align*}
\]

where \( \tilde{y} = \pi(y/W), \tilde{z} = \pi(z/W), \) and \( Z_{\text{res}}(\tilde{z}) = 1 + \frac{1}{2} \Gamma_{\text{bl}}(1 + i \tilde{z})^2 \). Note that at resonance, the horizontal velocity component is amplified by a factor \( \Gamma_{\text{bl}}^{-1} \) relative to the
wall velocity, \( v_{1y}^{d, \text{res}} \sim \Gamma_{bl}^{-1} v_{1z}^{d, \text{res}} \sim \Gamma_{bl}^{-1} V_1^{d, \text{res}} \), while the horizontal component is not.

In Fig. 2, we compare an effective (“Eff”) pressure-acoustics simulation of \( p_1 \) solving Eqs. (20a) and (25), with a full pressure-velocity simulation of \( p_1 \) and \( v_1 \) from Eq. (17) as in Muller and Bruus. The analytical results (“Ana”) for \( p_1^{\text{res}}, v_{1y}^{d, \text{res}} \), and \( v_{1z}^{d, \text{res}} \) in Eq. (69) are also plotted along the line \( y = \frac{1}{2} W \) in Figs. 2(a2), 2(b2), and 2(c2), respectively. The relative deviation between the full and effective fields outside the boundary layer are less than 0.1% even though the latter was obtained using only 5000 degrees of freedom (DoF) on the coarse mesh compared to the 600 000 DoF on the fine mesh for the former. The effective model \( v_{1y}^{d, \text{Eff}} \) gives the boundary-layer velocity \( v_{1y}^d \) by Eq. (22), and thus by Eq. (18) the complete velocity \( v_1 = v_{1y}^{d, \text{Eff}} + v_1^d \) (blue dots in Fig. 2).

To study the resonance behaviour of the acoustic resonator further, we compute the space- and time-averaged energy density \( \langle E_{ac}^{\text{d}} \rangle \) stored in the acoustic field for frequencies \( f \) close to the resonance frequency \( f_{res} \). Inserting \( k_y = \left( \frac{\pi}{W} \right) (1 + (i/2) \Gamma_{bl}) + \left( \frac{2 \pi}{c_0} (f - f_{res}) \right) \) into Eq. (65), results in the Lorentzian line-shape for \( \langle E_{ac}^{\text{d}} \rangle \),

\[
\langle E_{ac}^{\text{d}} \rangle = \frac{1}{4} \rho_0 \left( \frac{4 V_1^{d, \text{res}}}{\pi \Gamma_{bl}} \right)^2 \quad \text{for} \quad f \approx f_{res}.
\]

As shown in the graph of \( \langle E_{ac}^{\text{d}, \text{res}} \rangle \) in Fig. 3, there is full agreement between the effective pressure-acoustics model, the full pressure-velocity model, and the analytical model. In this figure we also show the result obtained using classical pressure acoustics (CPA, gray curves) with \( \delta_p p_1 = i \omega p_0 V_1^{d, \text{res}} \) where we see that the boundary layer introduces both damping and shift of the resonance frequency. From the resonance curve follows the maximum energy density at resonance, \( \langle E_{ac}^{\text{d}, \text{res}} \rangle = \langle E_{ac}^{\text{d}, (f_{res})} \rangle \), and the quality factor \( Q \),

\[
\langle E_{ac}^{\text{d}, \text{res}} \rangle = \frac{1}{4} \rho_0 \left( \frac{4 V_1^{d, \text{res}}}{\pi \Gamma_{bl}} \right)^2, \quad Q = \frac{1}{\Gamma_{bl}} = \frac{H}{\delta}.
\]

This is also in agreement with the Q-factor in Eq. (33),

\[
Q = \frac{2 \int_{W/2} \frac{1}{4} \delta_p p_1 |v_1^{d, \text{res}}|^2 \, dy \, dz}{2 \int_{W/2} \frac{1}{4} \rho_0 |v_1^{d, \text{res}}|^2 \, dy} = \frac{H}{\delta},
\]

which was previously derived by Muller and Bruus and by Hahn et al. using the approximation \( P_{\text{loss}} \approx P_{\text{diss}} \) in Eq. (33).

---

**FIG. 2.** (Color online) First-order pressure and velocity fields in the vertical rectangular cross section of a long, straight channel of width \( W = 380 \mu m \) and height \( H = 160 \mu m \) at resonance \( f_{res} = 1.967 \) MHz and actuation velocity \( V_1^{d, \text{res}} = 2 \eta_{\text{eff}} \times 0.1 \) nm. Color plots of the full (upper half) and effective (lower half) model fields: (a1) the pressure \( p_1 \) from \(-1 \) MPa (dark purple) to \( 1 \) MPa (light gray dashed lines) and the finite element mesh (gray), (b1) the horizontal velocity \( v_1 \) from \( 0 \) m/s (black) to \( 7 \) m/s (white), and (c1) the vertical velocity \( v_{1z} \) from \(-1 \) mm/s (black) to \( 1 \) mm/s (white). Line plots [Full, Eff, and Ana = analytics from Eq. (69)] at \( y_0 = \frac{1}{2} W \) for \(-\frac{1}{2} H < z < -\frac{1}{2} H + \frac{7}{2} \) (light gray dashed lines) of (a2) the relative pressure deviation \( p_1(y_0, z)/p_1(y_0, 0) - 1 \), (b2) the horizontal velocity \( v_{1y} \), and (c2) the vertical velocity \( v_{1z} \). Dots are the full velocity (18) \( v_1 = v_1^d + v_1^l \) with \( v_1^l \) from Eq. (22) and \( v_1^d \) from either “Eff” (dark blue dots) or “Ana” (light red dots). The insets are the corresponding plots along the entire line \(-\frac{1}{2} H < z < \frac{1}{2} H \).
B. Second-order streaming solution

For the full model at resonance $f_{\text{res}}$, we solve Eq. (34), while for the effective model we solve Eq. (52) with the boundary condition on $v_y^d$ obtained by inserting the velocity fields from Eq. (69) into Eq. (54). At the surfaces $z = \pm \frac{1}{2}H$, we find to lowest order in $\varepsilon$,

$$v_{2y}^0 = \frac{3}{8c_0} \left( \frac{4V_0}{\pi\Gamma_0} \right)^2 \sin(2\gamma),$$  \hspace{1cm} (73a)

$$v_{2y}^0 = -\left( k_0 \delta \right)^{1/2} \left( \frac{4V_0}{\pi\Gamma_0} \right)^2 \left[ 1 + 10 \cos(2\gamma) \right].$$  \hspace{1cm} (73b)

The resulting fields of the two models are shown in Fig. 4. Again, we have good quantitative agreement between the two numerical models, now better than 1% or 3$\varepsilon$ for 9000 DoF and 600 000 DoF, respectively.

Analytically, Eq. (73a) is the usual parallel-direction boundary condition for the classical Rayleigh streaming, while Eq. (73b) is beyond that, being the perpendicular-direction boundary condition on the streaming, which is a factor $k_0 \delta \approx 3 \times 10^{-3}$ smaller than the parallel one. This is confirmed in Fig. 4(b) showing the streaming velocity close to $z = -\frac{1}{2}H$ at $y = \frac{1}{4}W$.

IX. EXAMPLE II: A CURVED OSCILLATING CAVITY

Next, we implement in COMSOL our boundary conditions (25) and (55) in a system with a curved solid-fluid interface that oscillates in any direction, as described in Sec. VII. We consider an ellipsoidal fluid domain (water) of horizontal major axis $W = 380 \mu$m and vertical minor axis $H = 160 \mu$m surrounded by a rectangular solid domain (Pyrex) of width $W_{\text{sl}} = 680 \mu$m and height $H_{\text{sl}} = 460 \mu$m, see Fig. 5. We actuate the solid at its bottom surface using a vertical velocity amplitude $V_{2y}^{\text{act}} = d_0 \omega \sin(\pi y/W_{\text{sl}})$ with $d_0 = 0.1 \text{ nm}$ and at the resonance frequency $f_{\text{res}} = 2.222 \text{ MHz}$, which has
been determined numerically as in Fig. 3. The linear governing equations for the displacement field $u_d$ of the solid are those used by Ley and Bruus,\textsuperscript{48}

$$V \cdot \sigma_{dd} = -\rho_d c_t^2 (1 + i\Gamma_d) u_d, \quad \text{solid domain},$$

(74a)

$$-i\omega u_{dd} = V_{act}^2(y) e_z, \quad \text{actuation at } z = -\frac{1}{2} H_d,$$

(74b)

$$n_d \cdot \sigma_{dd} = 0, \quad \text{at solid–air interfaces},$$

(74c)

$$n_d \cdot \sigma_{dd} = n_d \cdot \sigma_1, \quad \text{at solid–fluid interfaces},$$

(74d)

where $\sigma_{dd} = \rho_d c_t^2 [V u_{dd} + (V u_{dd})^T] + \rho_d (c_t^2 - 2c_r^2) [V \cdot u_d]$ is the stress tensor of the solid with mass density $\rho_d$, transverse velocity $c_t$, longitudinal velocity $c_\alpha$, and damping coefficient $\Gamma_d$, while $n_d$ is the solid surface normal, and $n_d \cdot \sigma_1 = e_z \cdot \sigma_1$ is the fluid stress on the solid, Eq. (26). The material parameter values are listed in Table I.

We solve numerically Eqs. (20a) and (25) in first order and Eqs. (52) and (55) in second order. The results are shown in Fig. 5, where we compare the simulation results from the full boundary-layer resolved simulation of Eq. (34) with the effective model. Even for this more complex and realistic system consisting of an elastic solid with a curved oscillating interface coupled to a viscous fluid, we obtain good quantitative agreement between the two numerical models, better than 1% for 600 000 DoF and 9000 DoF, respectively.

X. CONCLUSION

We have studied acoustic pressure and streaming in curved elastic cavities having time-harmonic wall oscillations in any direction. Our analysis relies on the condition that both the surface curvature and wall displacement are sufficiently small as quantified in Eq. (16).

We have developed an extension of the conventional theory of first-order pressure acoustics by including the viscous effects of the thin viscous boundary layer. On this theory, we have also derived a slip-velocity boundary condition for the steady second-order acoustic streaming, which allows for efficient computations of the resulting incompressible Stokes flow.

The core of our theory is the decomposition of the first- and second-order fields into long- and short-range fields varying on the large bulk length scale $d$ and the small boundary-layer length scale $\delta$, respectively, see Eqs. (19) and (35). In the physically relevant limits, this velocity decomposition allows for analytical solutions of the boundary-layer fields. We emphasize that in contrast to the conventional second-order matching theory of inner solutions in the boundary layer and outer solutions in the bulk, our long- and short-range, second-order, time-averaged fields co-exist in the boundary layer, but the latter die out exponentially beyond the boundary layer leaving only the former in the bulk.

The main theoretical results of the extended pressure acoustics in Sec. III are the boundary conditions (25) and (26) for the pressure $p_1$ and the stress $\sigma_1 \cdot e_z$ expressed in terms of the pressure $p_1$ and the velocity $V_1^0$ of the wall. These boundary conditions are to be applied to the governing Helmholtz equation (20a) for $p_1$, and the gradient form (19b) of the compressional acoustic velocity field $v_1^0$. Furthermore, in Sec. IV, we have used the extended pressure boundary condition to derive an expression for the acoustic power loss $P_{loss}$, Eq. (32), and the quality factor $Q$, Eq. (33), for acoustic resonances in terms of boundary-layer and bulk loss mechanisms. The main results of the streaming theory in Sec. V are the governing incompressible Stokes equation (52) for the streaming velocity $v_2^0$ and the corresponding extended boundary condition (55) for the streaming slip velocity $v_2^0$. In this context, we have developed a compact formalism based on the $I^{(ab)}$-integrals of Eq. (43) to carry out with relative ease the integrations that lead to the analytical expression for $v_2^0$. Last, in Sec. VI, we have applied our extended pressure-acoustics theory to several special cases. We have shown how it leads to predictions that goes beyond previous theoretical results in the literature by Lord Rayleigh,\textsuperscript{22} Nyborg,\textsuperscript{33} Lee and Wang,\textsuperscript{34} and Vanneste and Bühler,\textsuperscript{35} while it does agree in the appropriate limits with these results.

The physical interpretation of our extended pressure acoustics theory may be summarized as follows: The fluid velocity $v_1$ is the sum of a compressible velocity $v_1^0$ and an incompressible velocity $v_1^1$, where the latter dies out beyond the boundary layer. In general, the tangential component $V_{1\|} = v_{1\|}^0 + v_{1\|}^1$ of the no-slip condition at the wall induces a tangential compression of $v_1^0$ due to the tangential compression of $v_1^0$ and $V_{1\|}^0$. This in turn induces a perpendicular velocity component $v_{1\perp}^0$ due to the incompressibility of $v_1^0$. To fulfill the perpendicular no-slip condition $V_{1\perp}^0 = v_{1\perp}^0 + v_{1\perp}^1$, the perpendicular component $v_{1\perp}^0$ of the acoustic velocity must therefore match not just the wall velocity $V_{1\perp}^0$, as in classical pressure acoustics, but the velocity difference $V_{1\perp}^0 - v_{1\perp}^0$. The inclusion of $v_{1\perp}^0$ takes into account the power delivered to the acoustic fields by the tangential wall motion, and the power lost from the acoustic fields due to tangential fluid motion. Consequently, by incorporating into the boundary condition an analytical solution of $v_1^0$, our theory leads to the correct acoustic fields, resonance frequencies, resonance $Q$-factors, and acoustic streaming.

In Secs. VII–IX we have demonstrated the implementation of our extended pressure acoustics theory in numerical finite-element COMSOL models, and we have presented the results of two specific models in 2D: a water domain with a rectangular cross section and a given velocity actuation on the domain boundary, and a water domain with an elliptic cross section embedded in a rectangular glass domain that is actuated on the outer boundary. By restricting our examples to 2D, we have been able to perform direct numerical simulations of the full boundary-layer-resolved model, and to use these results for validation of our extended acoustic pressure and streaming theory. Remarkably, we have found that even in 2D, our approach makes it possible to simulate acousto-fluidic systems with a drastic nearly 100-fold reduction in the necessary degrees of freedom, while achieving the same quantitative accuracy, typically of order $k_0 \delta$, compared to direct numerical simulations of the full boundary-layer resolved model. We have identified three reasons for this
reduction: (1) Neither our first-order nor our second-order method involve the fine-mesh resolution of the boundary layer. (2) Our first-order equations (20a) and (25) requires only the scalar pressure $p_1$ as an independent variable, while the vector velocity $v_1$ is subsequently computed from $p_1$, Eq. (19b). (3) Our second-order equations (52) and (55) avoid the numerically demanding evaluation in the entire fluid domain of large terms that nearly cancel, and therefore our method requires a coarser mesh compared to the full model, also in the bulk.

The results from the numerical examples in Secs. VIII and IX show that the extended pressure acoustics theory has the potential of becoming a versatile and very useful tool in the field of acoustofluidics. For the fluid-only rectangular domain in Sec. VIII, we showed how the theory not only leads to accurate numerical results for the acoustic fields and streaming, but also allows for analytical solutions, which correctly predict crucial details related to viscosity of the first-order acoustic resonance, and which open up for a deeper analysis of the physical mechanisms that lead to acoustic streaming. For the coupled fluid-solid system in 2D of an elliptical water domain embedded in a rectangular glass block, we showed in Sec. IX an important example of a more complete and realistic model of an actuated acoustofluidic system. The extended pressure acoustics theory allowed for calculations of acoustic fields and streaming with a relative accuracy lower than 1%. Based on preliminary work in progress in our group, it appears that the extended pressure acoustic theory makes 3D simulations feasible within reasonable memory consumptions for a wide range of microscale acoustofluidic systems such as fluid-filled cavities and channels driven by attached piezoelectric crystals as well as droplets in two-phase systems and on vibrating substrates.

Currently, we have neglected thermal dissipation. It would of course be an obvious and interesting study, to extend the presented theory to include thermoviscous effects. Previous studies on acoustofluidic systems with flat walls oscillating only in the perpendicular direction, have shown that the acoustic streaming is unaffected for channels with a height larger than 250 $\delta \approx 100 \mu$m. We thus expect the predictions of this work to hold for such “high” channels. However, for more flat channels, a significant reduction of the acoustic streaming is predicted, for example a reduction factor of 2 for a channel of height 25$\delta \approx 10 \mu$m. For such “flat” channels, the thermal boundary-layer must be included in our model to ensure reliable predictions. Such an extension of our model seems feasible, as the thermal boundary-layer width is about 3 times smaller than the viscous boundary-layer width for water at 2 MHz and 25 $^\circ$C. Thus the basic idea of a weakly curved, thin boundary-layer model can be maintained, but of course at the expense of analytical complications arising from including the heat transport equation together with temperature dependence of the material parameters in the presented first- and second-order perturbation theory.

Although we have developed the extended pressure-acoustics theory and corresponding streaming theory within the narrow scope of microscale acoustofluidics, our theories are of general nature and may likely find a much wider use in other branches of acoustics.

APPENDIX A: DIFFERENTIAL GEOMETRY

In the following we present the basic differential geometry used in this work. Because our analysis is carried out in the limit of weakly curved, thin boundary layers, defined by $\varepsilon \ll 1$ of Eq. (10) as discussed in Sec. II D, simplifications arise so that we do not need to unfold the full notation of differential geometry based on co- and contravariant derivatives, the metric tensor, and the full Christoffel symbols. Instead, we follow the tradition in the field set by Nyborg and by Lee and Wang, and use the vectorial notation based on the unit tangent vectors $e_i$ and the scale factors $h_i$ at position $r$ in the thin boundary layer,

$$e_i = \frac{1}{h_i} \partial_i r, \quad \text{with} \quad h_i = |\partial_i r| \quad \text{for} \quad i = \xi, \eta, \zeta, \quad (A1a)$$

$$e_k \cdot e_i = \delta_{ki}, \quad \text{orthonormality by construction.} \quad (A1b)$$

It is natural to introduce the scaled derivatives $\tilde{\partial}_i$ and the curvilinear quantities $T_{\tilde{k}ji}$ and $\mathcal{H}_{\tilde{k}}$,

$$\tilde{\partial}_j = h_j^{-1} \partial_j, \quad \text{so that} \quad e_i = \tilde{\partial}_i r, \quad (A2a)$$

$$T_{\tilde{k}ji} = (\tilde{\partial}_j e_i) \cdot e_i, \quad \text{for} \quad k, j, i = \xi, \eta, \zeta, \quad (A2b)$$

$$\mathcal{H}_{\tilde{k}} = T_{\tilde{k}ii}, \quad \text{sum over repeated index} \quad i. \quad (A2c)$$

$T_{\tilde{k}ji}$ is related to, but not identical with, the celebrated Christoffel symbols of differential geometry. The following relations for $T_{\tilde{k}ji}$ are useful in the analysis:

$$T_{\tilde{k}ji} = T_{ij\tilde{k}}, \quad \text{for} \quad i \neq j, \quad (A3a)$$

$$T_{\tilde{k}\xi \xi} = T_{\eta \eta \xi} = T_{\zeta \zeta \xi} = 0, \quad (A3b)$$

$$T_{\tilde{j}i} = 0. \quad (A3c)$$

Equation (A3a) follows from $T_{\tilde{k}ji} = ((1/h_{\tilde{k}}/h_j)(1/h_i)\partial_j(r)) \cdot e_i = ((\partial_j \partial_i r/h_i h_j) \cdot e_i) - (\partial_j h_i h_j) e_i \cdot e_i$, which is symmetric in $k$ and $j$ as the last term is zero for $j \neq i$, Eq. (A3b) is proven by observing $T_{\tilde{k}\xi \xi} = (\tilde{\partial}_\xi e_\xi) \cdot e_\xi = \frac{1}{2} \tilde{\partial}_k (e_\zeta \cdot e_\zeta) = 0$ as $e_\zeta \cdot e_\zeta = 1$, and Eq. (A3c) arises because $\zeta$ is defined as the normal direction, and as a consequence $T_{\tilde{k}ij}$ is only nonzero for the tangential derivatives $T_{\tilde{j}ij}$ and $T_{\eta ij}$, and $\mathcal{H}_{\tilde{k}} = T_{\tilde{k}\xi \xi}$. It is in this sense that the surface length scale is set by $R \sim \min \{T_{\tilde{k}ji}, \mathcal{H}^{-1}_{\tilde{k}}\}$ as stated in Sec. II E.

From now on, we use the index notation, where, as in Eq. (A2c), a repeated index implies a summation. In curvilinear coordinates, the $\textbf{V}$ operator and vector fields are written as

$$\textbf{V} = e_i \tilde{\partial}_i, \quad \textbf{A} = A_i e_i, \quad (A4)$$

and from this all other differential operators are calculated. The first example is the Laplacian of a scalar,

$$\nabla^2 g = \nabla \cdot \nabla g = (e_i \tilde{\partial}_i) \cdot e_i \tilde{\partial}_i g$$

$$= [e_i \cdot e_i \tilde{\partial}_i + e_j \cdot (\tilde{\partial}_j e_i)] \tilde{\partial}_i g$$

$$= (\tilde{\partial}_i + T_{\tilde{j}i}) \tilde{\partial}_i g = (\tilde{\partial}_i \tilde{\partial}_i + \mathcal{H}_{\tilde{k}} \tilde{\partial}_k) g. \quad (A5)$$
The divergence of a vector field $A$ takes the form
\[ \nabla \cdot A = e_i \cdot \partial_i (A_k e_k) = \left( \partial_i A_k \right) e_k + A_k \partial_i e_k \cdot e_i = (\partial_i A_k) \delta_{ik} + A_k T_{ik} = (\partial_i + H_k) A_k, \] (A6)
while the gradient of a vector field $B$ is
\[ \nabla B = e_i \partial_i (B_k e_k) = e_k \left( \partial_i B_k \right) e_i + B_j \partial_j e_i = e_k (\partial_i B_i + T_{ik} B_j) e_i. \] (A7)

From this follows the advective derivative $(A \cdot \nabla)B$ of a vector $B$ with respect to a vector $A$,
\[ (A \cdot \nabla)B = A_k (\partial_i B_i + T_{ik} B_j) e_i, \] (A8)
and the Laplacian of a vector $B$,
\[ \nabla^2 B = \nabla \cdot (A \cdot \nabla)B = e_n \partial_n \left[ e_k (\partial_i B_i + T_{ik} B_j) e_i \right] = H_k (\partial_i B_i + T_{ik} B_j) e_i + (\partial_i B_i + T_{ik} B_j) \partial_k e_i \] (A9)

In the analysis of fields in the weakly curved, thin boundary layer, it is useful to decompose a given vector $A$ into parallel and perpendicular components,
\[ A = A_\parallel + A_\perp e_\perp, \] (A10a)
\[ A_\parallel = A_\zeta e_\zeta + A_\eta e_\eta = A_\zeta e_\zeta, \] (A10b)
where here and in the following, repeated Greek index $\zeta$ only sums over the tangential indices $\zeta$ and $\eta$. Likewise, the parallel components of the $\nabla$ operator (A4), the divergence (A6), and the advective derivative (A8) are
\[ \nabla_\parallel = e_\zeta \partial_\zeta, \] (A11a)
\[ \nabla_\parallel \cdot A_\parallel = (\partial_\zeta + H_\zeta) A_\zeta, \] (A11b)
\[ (A \cdot \nabla_\parallel)B = A_\zeta (\partial_\zeta B_\zeta + T_{\zeta\eta} B_\eta) e_\zeta, \] (A11c)

For the short-ranged boundary-layer vector field $A^\delta = A^\delta(\xi, \eta) a^\delta(\zeta)$ introduced in Eq. (14), and the analogous scalar field $g^\delta = g^{\delta(\xi, \eta) a^\delta(\zeta)}$, the derivative expressions simplifies in the weakly curved thin boundary-layer limit $\epsilon \ll 1$. The reason is that terms containing surface-derivative quantities $V_\parallel$, $T_{\parallel\perp}$, and $H_k$, all of size $d_1^{-1}$, are a factor of $\epsilon$ smaller than terms with the parallel derivative $\partial_\zeta = \partial_\zeta$, which picks up a factor of $\delta^{-1}$ due to the factor $a^\delta(\zeta)$ that decays on the length scale $\delta$.

For $\nabla^2 g^\delta$ in Eq. (A5), the first term $\partial_i \partial_\zeta g^\delta \sim \partial_\zeta^2 g^\delta$ is $\epsilon^{-1}$ larger than the second term $H_k \partial_\zeta g^\delta \sim R^{-1} \delta^{-1} g^\delta$, so that $\nabla^2 g^\delta \sim \partial_\zeta^2 g^\delta$ as stated in Eq. (15a).

Similarly, for $\nabla^2 A^\delta$ in Eq. (A9), the only term that does not contain at least one factor $H_k$ or $T_{\parallel\perp}$ is $(\partial_k \partial_\zeta A^\delta_k) e_i \approx (\partial_\zeta^2 A^\delta_k) e_i = \partial_\zeta^2 A^\delta$ as stated in Eq. (15b).

Finally, for Eq. (A6), $\nabla \cdot A^\delta = \nabla_\parallel \cdot A^\delta + (\partial_\zeta + H_\zeta) A^\delta_\zeta \approx \nabla_\parallel \cdot A^\delta + \partial_\zeta A^\delta_\zeta$ as stated in Eq. (15c).

### APPENDIX B: ACOUSTIC POWER BALANCE

The time averages $\langle E_{\text{kin}}^{\text{ac}}, E_{\text{pot}}^{\text{ac}}, \rangle$ and $\langle E_{\text{ac}} \rangle$ of the kinetic, the potential, and the total acoustic energy density, respectively, are given by
\[ \langle E_{\text{kin}}^{\text{ac}} \rangle = \frac{1}{2} \rho_0 \langle v_1 \cdot v_1 \rangle, \] (B1a)
\[ \langle E_{\text{pot}}^{\text{ac}} \rangle = \frac{1}{2} \kappa_0 \langle p_1 p_1 \rangle, \] (B1b)
\[ \langle E_{\text{ac}} \rangle = \langle E_{\text{kin}}^{\text{ac}} \rangle + \langle E_{\text{pot}}^{\text{ac}} \rangle. \] (B1c)

Using Gauss’s theorem and $\rho_0 \partial_\zeta v_1 = V \cdot \sigma_1$, the time-averaged total power delivered by the surrounding wall is written as the sum of the time-averaged rate of change of the acoustic energy and total power dissipated into heat,
\[ \int_{\partial \Omega} \langle V_0^0 \cdot \sigma_1 \rangle \cdot n \, dA = \int_{\Omega} V \cdot \langle v_1 \cdot \sigma_1 \rangle \, dV \] (B2a)
\[ = \int_{\Omega} \langle |v_1 \cdot (V \cdot \sigma_1)| + (\langle V v_1 \rangle : \sigma_1) \rangle \, dV \] (B2b)
\[ = \int_{\Omega} \langle |\partial_\zeta E_{\text{ac}}| + (\langle V v_1 \rangle : \tau_1) \rangle \, dV. \] (B2c)

Solving for the time-averaged change in acoustic energy $\int_{\Omega} \langle \partial_\zeta E_{\text{ac}} \rangle \, dV$ in Eq. (B2c) gives
\[ \int_{\Omega} \langle \partial_\zeta E_{\text{ac}} \rangle \, dV = \int_{\partial \Omega} \langle V_0^0 \cdot \tau_1 \rangle \cdot n \, dA - \int_{\Omega} \langle (V v_1) : \tau_1 \rangle \, dV \] (B3a)
\[ = \int_{\partial \Omega} \langle V_0^0 (-p_1) \rangle \cdot n \, dA + \int_{\Omega} \langle v_1 \cdot (V \cdot \tau_1) \rangle \, dV, \] (B3b)
where Gauss’s theorem transforms $\int_{\partial \Omega} \langle V_0^0 \cdot \tau_1 \rangle \cdot n \, dA$ into a volume integral, and $n = -e_\zeta$ is the normal vector of the fluid domain $\Omega$. We may interpret the last term in Eq. (B3b) as the rate of change of stored energy due to the viscosity-induced power $\langle P_{\text{visc}} \rangle$.
\[ \langle P_{\text{visc}} \rangle = \langle P_{\text{diss visc}} \rangle + \langle P_{\text{wall visc}} \rangle. \] (B4)

Here, $\langle P_{\text{diss visc}} \rangle$ is the viscous power dissipation into heat and $\langle P_{\text{wall visc}} \rangle$ is the power from the viscous part of the work performed by the wall on the fluid,
\[ \langle P_{\text{visc}} \rangle = \int_{\Omega} \langle v_1 \cdot (V \cdot \tau_1) \rangle \, dV, \] (B5a)
\( \langle P_{\text{visc}} \rangle = \int_\Omega \langle (\mathbf{V} \cdot \mathbf{v}_1) : \mathbf{\tau}_1 \rangle \, dV \),  
(B5b)

\( \langle P_{\text{wall}} \rangle = \oint_{\partial \Omega} \langle \mathbf{v}_1 \cdot \mathbf{\tau}_1 \rangle \cdot n \, dA. \)  
(B5c)

Using Eqs. (18) and (19) we evaluate \( \langle P_{\text{visc}} \rangle \),

\[ \langle P_{\text{visc}} \rangle = \int_\Omega \left( \mathbf{v}_1 \cdot (i\Gamma p_1 - i\omega \rho_1 \mathbf{v}_1^\Gamma) \right) \, dV \]
(B6a)

\[ = \int_\Omega \left[ -\frac{\Gamma \omega \rho_1}{2} |\mathbf{v}_1|^2 + \langle \partial \mathbf{E}_{\infty}^{\text{kin},\mathbf{A}} \rangle \right] \, dV \]
(B6b)

\[ - \oint_{\partial \Omega} \langle p_1 \mathbf{v}_1^\Gamma \rangle \cdot n \, dA, \]
(B6c)

where we used Eq. (19) and Gauss’s theorem. Inserting Eq. (B6c) into Eq. (B3b) leads to Eq. (30). Comparing with Eq. (32), we can relate \( \langle P_{\text{loss}} \rangle = \langle P_{\text{diss}}^\Gamma \rangle \) and \( \langle P_{\text{visc}} \rangle \),

\[ \langle P_{\text{loss}} \rangle = \langle P_{\text{visc}} \rangle - \oint_{\partial \Omega} \left( \mathbf{p}_1 \left[ \frac{i}{k_s} \mathbf{V} \cdot \mathbf{V}_1^0 \right] \right) \cdot n \, dA \]  
(B7a)

\[ = \langle P_{\text{diss}}^\Gamma \rangle + \langle P_{\text{wall}} \rangle - \oint_{\partial \Omega} \left( \mathbf{p}_1 \left[ \frac{i}{k_s} \mathbf{V} \cdot \mathbf{V}_1^0 \right] \right) \cdot n \, dA. \]  
(B7b)

Note that \( \langle P_{\text{loss}} \rangle \) is not in general the same as the power \( \langle P_{\text{visc}} \rangle \) dissipated into heat. These might, however, be approximately equal if the power

\[ \oint_{\partial \Omega} \langle p_1 \mathbf{V}_1^\Gamma \rangle \cdot n \, dA \]

delivered by the pressure is approximately balanced by dissipation \( \langle P_{\text{diss}}^\Gamma \rangle \). This happens, if

\[ \oint_{\partial \Omega} \langle p_1 \mathbf{V}_1^\Gamma \rangle \cdot n \, dA \]

is much larger than \( \langle P_{\text{wall}} \rangle \) and

\[ \oint_{\partial \Omega} \left( \mathbf{p}_1 \left[ \frac{i}{k_s} \mathbf{V} \cdot \mathbf{V}_1^0 \right] \right) \cdot n \, dA, \]

which is usually satisfied.

**APPENDIX C: COMPARISON WITH PREVIOUS RESULTS IN THE LITERATURE**

1. **Comparison with Lee and Wang (1989)**

In the following we rewrite the \( \zeta \)-component \( v_{2 \zeta}^0 \) of the streaming velocity in Eq. (58a) using the notation of Lee and Wang\(^{34}\) to compare it directly with \( u_L \) given in their Eq. (19)\(^LW\). First, we take the complex conjugate of the argument of the real value, set \( \beta = \xi \), and write explicitly the sum over the repeated index \( z = \xi, \eta \).

\[ v_{2 \zeta}^0 = -\frac{1}{4\epsilon} \text{Re} \left\{ v_{1 \xi}^0 \left( \overline{\partial \xi v_{1 \eta}^0} + v_{1 \xi}^0 \overline{\partial \eta v_{1 \zeta}^0} \right) + v_{1 \eta}^0 \overline{\partial \eta v_{1 \zeta}^0} \right\} \]
\[ + v_{1 \eta}^0 \left( \overline{\partial v_{1 \xi}^0} + v_{1 \xi}^0 \overline{\partial v_{1 \eta}^0} \right) + \left( 2i + i \right) \overline{\mathbf{v} \cdot \mathbf{v}_1^0} \]
\[ - 2(2 + 3i) \partial_z \overline{v_{1 \xi}^0} - (2 - i) \overline{H_\xi v_{1 \eta}^0 v_{1 \zeta}^0} \]
\[ - 2i(2 + 3i) \partial_z \overline{v_{1 \eta}^0} + v_{1 \xi}^0 \overline{H_\eta v_{1 \zeta}^0 v_{1 \xi}^0} \}. \]  
(C1)

This expression is rewritten using the notation of Lee and Wang, \( v_{1 \eta}^0 = u_{a0} \), \( (v_{1 \xi}^0, v_{1 \eta}^0, v_{1 \zeta}^0) = (u_{a0}, \nu_{a0}, w_{a0}) \), \( \langle \xi, \eta, \zeta \rangle = (x, y, z) \), \( H_\xi = H_\xi(\overline{\partial \xi}, \overline{\partial \eta}, \overline{\partial \zeta}) = (\partial_{x1}, \partial_{y2}, \partial_{z3}) \), \( T_{ijk} = T_{ijkl}^LW \), and \( V_{1 \eta}^0 = w_{a0} + O(\epsilon) \).

\[ v_{2 \eta}^0 = -\frac{1}{4\epsilon} \text{Re} \left\{ u_{a0} \left[ \partial_z u_{a0}^* + v_{a0} T_{ijkl}^LW \right] \right\} \]
\[ + \nu_{a0} \left[ \partial_{x1} u_{a0}^* + v_{a0} T_{ij13}^LW + u_{a0} \left( 2 + i \right) \overline{\mathbf{v} \cdot \mathbf{u}_{a0}} \right. \]
\[ - 2(2 + 3i) \partial_{x1} w_{a0} - (2 - i) Hw_{a0} \]
\[ - 2i w_{a0}^* \left( \partial_z u_{a0} + u_{a0} T_{13i}^LW + v_{a0} T_{ij13}^LW \right) \]
\[ = -\frac{1}{4\epsilon} \text{Re} \left\{ u_{a0} \left[ \partial_z u_{a0}^* + v_{a0} T_{ijkl}^LW \right] \right\} \]
\[ - 2(2 + 3i) \partial_{x1} w_{a0} - (2 - i) Hw_{a0} \]
\[ - 2i w_{a0}^* \left( \partial_z u_{a0} + u_{a0} T_{13i}^LW + v_{a0} T_{ij13}^LW \right) \]
\[ = u_L + \frac{1}{2\epsilon} \text{Re} \left\{ iw_{a0}^* \left( \partial_z u_{a0} + u_{a0} T_{13i}^LW + v_{a0} T_{ij13}^LW \right) \right\}. \]  
(C2)

Here, we have used the definition of \( q_\eta \) in Eq. (3)\(^LW\) and the result in Eq. (19)\(^LW\) for the \( x \)-component \( u_L \) of the streaming velocity just outside the boundary layer.

Similarly, we obtain for the \( \eta \)-component \( v_{2 \eta}^0 \) in Eq. (58a),

\[ v_{2 \eta}^0 = v_L + \frac{1}{2\epsilon} \text{Re} \left\{ iw_{a0}^* \left( \partial_z u_{a0} + u_{a0} T_{13i}^LW + v_{a0} T_{ij13}^LW \right) \right\}, \]  
(C3)

where we have used the definition in Eq. (4)\(^LW\) for \( q_\eta \) and the result in Eq. (20)\(^LW\) for the \( y \)-component \( v_L \) of the streaming velocity. The comparison obtained in Eqs. (C2) and (C3) is discussed in Sec. VI A.

2. **Comparison with Vanneste and Bühler (2011)**

In the following we rewrite the \( \zeta \)-component \( v_{2 \zeta}^0 \) of the Eulerian streaming velocity in Eq. (59b) using the notation of Vanneste and Bühler\(^{35}\) to compare it directly with the Lagrangian streaming velocity \( \mathbf{u}_{\text{slip}}^L \) given in their Eq. (4.10)\(^BV\). First, given the flat wall, \( \langle \xi, \eta, \zeta \rangle = (x, y, z) \) and for all \( i, j, k \) we have \( h_i = 1 \), \( h_j = 0 \), and \( T_{ijk} = 0 \). Then, we identify our first-order velocity fields with theirs: From Eqs. (2.6)\(^BV\) and (3.1)\(^BV\) follows \( v_{\phi}^2 = 2 \overline{\partial \phi} \), and from Eqs. (3.9)\(^BV\) and (3.10)\(^BV\) we read that \( v_{\eta}^0 = -2 \overline{U}_e v_{\eta} - 2 \overline{V} e_{\eta} \). Next, we relate our steady Eulerian second-order velocity \( v_{\eta}^2 \) with their Lagrangian mean flow \( \mathbf{u}_{\text{slip}}^L \) Eq. (4.10)\(^BV\). Using \( \text{Re}Z = \frac{1}{2} \overline{Z} + c.c. = \frac{1}{2} Z^* + c.c. \), we obtain
\begin{equation}
\nu_{2g}^{0} = -\frac{1}{4\omega} \text{Re} \left\{ (1 - 2i) \nu_{1g}^{0} \partial_{x} \nu_{1g}^{0} - 4i \nu_{1g}^{0} \partial_{x} \nu_{1g}^{0} + \left[ (2 + i) \partial_{x} \nu_{1g}^{0} + 2i \left( \partial_{x} \partial_{x} v_{1g}^{0} - \partial_{y} \partial_{z} v_{1g}^{0} \right) \right] \nu_{1g}^{0} - 2i \nu_{1g}^{0} \partial_{x} \nu_{1g}^{0} \right\}
\end{equation}

\begin{equation}
= -\frac{1}{2\omega} \left\{ (1 + 2i) \nu_{0} \partial_{x} \nu_{0} + (1 + 2i) V \partial_{x} \nu_{0} + 4i \nu_{0} \partial_{x} \nu_{0} \partial_{x} \phi + 4i \nu_{0} \partial_{y} \partial_{z} \phi \right\} + \left( \partial_{x} \nu_{0} \partial_{x} \nu_{0} + (2 + i) \nu_{0} \partial_{x} \nu_{0} + 2i \left( \partial_{x} \nu_{0} + \partial_{x} \nu_{0} + \partial_{x} \phi \right) \right) \nu_{0} \right\} + c.c. - \frac{1}{\omega} (iv_{1g}^{1} \cdot v_{1g}^{1})
\end{equation}

\begin{equation}
= -\nu_{s}^{L} - \nu_{s}^{E} = \nu_{s}^{E}.
\end{equation}

Here, we have used that the Lagrangian velocity \( \nu_{s}^{E} \), calculated by Vanneste and Bühler in their Eq. (4.10)\(^1\), is related to the Eulerian velocity \( \nu_{s}^{E} \) through the Stokes drift velocity \( \nu_{s}^{E} = (1/\omega) \left( iv_{1g}^{1} \cdot v_{1g}^{1} \right) \) of Eq. (34c).

Similarly, for the \( y \) component \( \nu_{2g}^{0} \) of the streaming,

\begin{equation}
\nu_{2g}^{0} = \nu_{s}^{L} - \nu_{s}^{E} = \nu_{s}^{E}.
\end{equation}

The comparison obtained in Eqs. (C4) and (C5) is discussed in Sec. VI.B.


2T. Laurell and A. Lenshof, Microscale Acoustofluidics (Royal Society of Chemistry, Cambridge, 2015).


COMSOL Multiphysics 5.3a (2017), www.comsol.com (Last viewed 3 August 2018).
