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Theoretical study of time-dependent, ultrasound-induced acoustic streaming in microchannels

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Based on first- and second-order perturbation theory, we present a numerical study of the temporal buildup and decay of unsteady acoustic fields and acoustic streaming flows actuated by vibrating walls in the transverse cross-sectional plane of a long straight microchannel under adiabatic conditions and assuming temperature-independent material parameters. The unsteady streaming flow is obtained by averaging the time-dependent velocity field over one oscillation period, and as time increases, it is shown to converge towards the well-known steady time-averaged solution calculated in the frequency domain. Scaling analysis reveals that the acoustic resonance builds up much faster than the acoustic streaming, implying that the radiation force may dominate over the drag force from streaming even for small particles. However, our numerical time-dependent analysis indicates that pulsed actuation does not reduce streaming significantly due to its slow decay. Our analysis also shows that for an acoustic resonance with a quality factor \( Q \), the amplitude of the oscillating second-order velocity component is \( Q \) times larger than the usual second-order steady time-averaged velocity component. Consequently, the well-known criterion \( v_1 \ll c_s/Q \) for the validity of the perturbation expansion is replaced by the more restrictive criterion \( v_1 \ll c_s/Q \). Our numerical model is available as supplemental material in the form of COMSOL model files and MATLAB scripts.

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I. INTRODUCTION

Acoustophoresis has successfully been used in many applications to manipulate particles in the size range from about 0.5 mm down to about 2 \( \mu \)m [1]. However, for smaller particles, the focusing by the acoustic radiation force is hindered by the drag force from the suspending liquid, which is set in motion by the generation of an acoustic streaming flow [2,3]. This limits the use of acoustophoresis to manipulate submicrometer particles, relevant for application within medical, environmental, and food sciences, and it underlines a need for better understanding of acoustic streaming and ways to circumvent this limitation.

The phenomenon of acoustic streaming was first described theoretically by Lord Rayleigh [4] in 1884, and has later been revisited, among others, by Schlichting [5], Nyborg [6], Hamilton [7,8], Rednikov and Sadhal [9], and Muller et al. [10], to extend the fundamental treatment of the governing equations and to solve the equations for various open and closed geometries.

Numerical methods have been applied in many studies to predict the streaming phenomena observed in various experiments. Muller et al. [2] developed a numerical scheme to solve the acoustic streaming in the cross section of a long straight microchannel, which resolved the viscous acoustic boundary layers and described the interplay between the acoustic scattering force and the streaming-induced drag force on suspended particles. This scheme was later extended to take into account the thermoviscous effects arising from the dependence of the fluid viscosity on the oscillating temperature field [11]. Lei et al. [12,13] have developed a numerical scheme based on the effective slip-velocity equations, originally proposed by Nyborg in 1953 [14,15], which avoid the resolution of the thin boundary layers but still enable qualitative predictions of the three-dimensional streaming flows observed in microchannels and flat microfluidic chambers. To obtain quantitative results from such models that do not resolve the acoustic boundary layers, Hahn et al. [16] developed an effective model to determine the loss associated with the viscous stresses inside the thermoacoustic boundary layers, and apply this loss as an artificial bulk absorption coefficient. This enables the calculation of correct acoustic amplitudes, without resolving the thin acoustic boundary layers. Acoustic streaming in the cross section of a straight PDMS microchannel excited by surface acoustic waves was studied numerically by Nama et al. [17], describing the influence of the acoustically soft PDMS wall on the particle focussability, and examining the possibilities of having two tunable counterpropagating surface acoustic waves.

All of the above-mentioned studies consider steady acoustic streaming flows. This is reasonable as the streaming flow reaches steady state typically in a few milliseconds, much faster than other relevant experimental time scales. Furthermore, this allows for analytical solutions for the streaming velocity field in some special cases, and it makes it much easier to obtain numerical solutions. However, an experimental study by Hoyos and Castro [18] indicates that a pulsed actuation, instead of steady, can reduce the drag force from the streaming flow relative to the radiation force and thus allow the latter also to dominate manipulation of submicrometer particles. This might provide an alternative method to the one proposed by Antfolk et al. [19], which used an almost square channel with overlapping resonances to create a streaming flow that did not counteract the focusing of submicrometer particles. To theoretically study the effects of a pulsed ultrasound actuation, we need to solve the temporal evolution of the acoustic resonance and streaming, which is the topic of the present work. Numerical solutions of the time-domain acoustic equations were used by Wang and Dual [20] to calculate the time-averaged radiation force on a cylinder and the steady streaming...
around a cylinder, both in a steady oscillating acoustic field. However, they did not present an analysis of the unsteady buildup of the acoustic resonance and the streaming flow.

In this paper, we derive the second-order perturbation expansion of the time-dependent governing equations for the acoustic fields and streaming velocity, and solve them numerically for a long straight channel with acoustically hard walls and a rectangular cross section. The analysis and results are divided into two sections: (1) A study of the transient buildup of the acoustic resonance and streaming from an initially quiescent state towards a steady oscillating acoustic field and a steady streaming flow. (2) An analysis of the response of the acoustic field and the streaming flow to pulsed actuation, and quantifying whether this can lead to better focussability of submicrometer particles.

In previous studies, such as [2,11,17], only the periodic state of the acoustic resonance and the steady time-averaged streaming velocity are solved. When solving the time-dependent equations, we obtain a transient solution, which may also be averaged over one oscillation period to obtain an unsteady time-averaged solution.

II. BASIC ADIABATIC ACOUSTIC THEORY

In this section we derive the governing equations for the first- and second-order perturbations to unsteady acoustic fields in a compressible Newtonian fluid. We only consider acoustic perturbation in fluids, and treat the surrounding solid material as ideal rigid walls, a good approximation for water channels in glass-silicon systems. Moreover, extensive theoretical and experimental work [2,3,10,11,19,21] has shown that the adiabatic model describes the observed phenomena qualitatively correctly, while thermoviscous effects may lead to relative quantitative changes up to 30%. Thus by employing an adiabatic model, we can provide an analysis of experimental relevance while at the same time restricting the complexity and vast parameter space of the full problem. We leave a more complete thermoviscous analysis for future work. Our treatment is based on textbook adiabatic acoustics [22] and our previous study Ref. [11] of the purely periodic state.

A. Adiabatic thermodynamics

We employ the adiabatic approximation, which assumes that the entropy is conserved for any small fluid volume [23]. Consequently, the thermodynamic state of the fluid is described by only one independent thermodynamic variable, which we choose to be the pressure \( p \). See Table I for parameter values. The changes \( d\rho \) in the density \( \rho \) from the equilibrium state are given by

\[
\frac{d\rho}{\rho} = \kappa_s \frac{dp}{p},
\]

where the isentropic compressibility \( \kappa_s \) is defined as

\[
\kappa_s = \frac{1}{\rho} \left( \frac{\partial \rho}{\partial p} \right)_g = \frac{1}{\rho c_s^2}.
\]

B. Governing equations

Mass conservation implies that the rate of change \( \partial_t \rho \) of the density in a test volume with surface normal vector \( n \) is given by the influx (direction \( -n \)) of the mass current density \( \rho \mathbf{v} \). In differential form by Gauss’s theorem it is

\[
\partial_t \rho = \nabla \cdot [ -\rho \mathbf{v} ].
\]

Substituting \( \partial_t \rho \) and \( \nabla \rho \) using Eq. (1), and dividing by \( \rho \), the continuity equation (3a) becomes

\[
\kappa_s \partial_t \rho = - \nabla \cdot \mathbf{v} - \kappa_s \mathbf{v} \cdot \nabla \rho.
\]

Similarly, momentum conservation implies that the rate of change \( \partial_t (\rho \mathbf{v}) \) of the momentum density in the same test volume is given by the stress forces \( \mathbf{\sigma} \) acting on the surface (with normal \( n \)), and the influx (direction \( -n \)) of the momentum current density \( \rho \mathbf{v} \mathbf{v} \). In differential form, neglecting body forces, this becomes

\[
\partial_t (\rho \mathbf{v}) = \nabla \cdot [ \mathbf{\tau} - \rho \mathbf{v} \mathbf{v} ],
\]

where the viscous stress tensor is defined as

\[
\mathbf{\tau} = \eta [ \nabla \mathbf{v} + (\nabla \mathbf{v})^\top ] + [ \eta^b - \frac{2}{3} \eta ] (\nabla \cdot \mathbf{v}) \mathbf{I}.
\]

Here \( \mathbf{I} \) is the unit tensor and the superscript “T” indicates tensor transposition. Using the continuity equation (3a), the momentum equation (4a) is rewritten into the well-known Navier-Stokes form,

\[
\rho \partial_t \mathbf{v} = \nabla \cdot [ \mathbf{\tau} - \rho \mathbf{v} \mathbf{v} ],
\]

which is useful when solving problems in the time domain. Equations (3b) and (4c) constitute the nonlinear governing equations which we will study by applying the usual perturbation approach of small acoustic amplitudes.

C. First-order time-domain equations

The homogeneous, isotropic, quiescent thermodynamic equilibrium state is taken to be the zeroth-order state in the acoustic perturbation expansion. Following standard perturbation theory, all fields \( g \) are written in the form \( g = g_0 + g_1 \), for which \( g_0 \) is the value of the zeroth-order state, and \( g_1 \) is the acoustic perturbation which by definition has to be much smaller than \( g_0 \). For the velocity, the value of the zeroth-order state is \( \mathbf{v}_0 = \mathbf{0} \), and thus \( \mathbf{v} = \mathbf{v}_1 \). The zeroth-order terms solve the governing equations in the zeroth-order state and thus drop out of the equations. Keeping only first-order terms, we obtain the following first-order equations.

The first-order continuity equation (3b) becomes

\[
\kappa_s \partial_t \rho_1 = - \nabla \cdot \mathbf{v}_1,
\]

\[
\begin{array}{|c|c|c|c|}
\hline
\text{Parameter} & \text{Symbol} & \text{Value} & \text{Unit} \\
\hline
\text{Acoustic properties:} & & & \\
\text{Mass density} & \rho_0 & 9.971 \times 10^2 & \text{kg m}^{-3} \\
\text{Speed of sound} & c_s & 1.497 \times 10^3 & \text{m s}^{-1} \\
\text{Compressibility} & \kappa_s & 4.477 \times 10^{-10} & \text{Pa}^{-1} \\
\text{Transport properties:} & & & \\
\text{Shear viscosity} & \eta & 8.900 \times 10^{-4} & \text{Pa s} \\
\text{Bulk viscosity} & \eta^b & 2.485 \times 10^{-3} & \text{Pa s} \\
\hline
\end{array}
\]
and likewise, the momentum equation (4c) becomes

\[ \rho_0 \partial_t \mathbf{v}_1 = \nabla \cdot [\mathbf{r}_1 - \mathbf{p}_1 \mathbf{I}] \]  

(6a)

where \( \mathbf{r}_1 \) is given by

\[ \mathbf{r}_1 = \eta_0 [\nabla \mathbf{v}_1 + (\nabla \mathbf{v}_1)^T] + \left[ \eta_0^b - \frac{\zeta}{2} \eta_0 \right] (\nabla \cdot \mathbf{v}_1) \mathbf{I}. \]  

(6b)

Equations (5) and (6) determine together with a set of boundary conditions the time evolution of the first-order acoustic fields \( \mathbf{p}_1 \) and \( \mathbf{v}_1 \).

D. Second-order time-domain equations

Moving on to second-order perturbation theory, we write the fields as \( \mathbf{r} = \mathbf{r}_0 + \mathbf{r}_1 + \mathbf{r}_2 \), with \( \mathbf{r}_0 \) and \( \mathbf{r}_2 \) depending on both time and space. For simplicity and in contrast to Ref. [11], we do not include perturbations in \( \eta \) and \( \eta^b \). This will cause the magnitude of the streaming to be slightly off, as does the adiabatic approximation; however the qualitative behavior is not expected to change. The second-order time-domain continuity equation (3b) becomes

\[ \kappa_s \partial_t \mathbf{p}_2 = -\nabla \cdot \mathbf{v}_2 - \kappa_s \mathbf{v}_1 \cdot \nabla \mathbf{v}_1, \]  

(7)

and the momentum equation (4c) takes the form

\[ \rho_0 \partial_t \mathbf{v}_2 = -\rho_1 \partial_t \mathbf{v}_1 + \nabla \cdot [\mathbf{r}_2 - \mathbf{p}_2 \mathbf{I}] - \rho_0 (\mathbf{v}_1 \cdot \nabla) \mathbf{v}_1, \]  

(8a)

where \( \mathbf{r}_2 \) is given by

\[ \mathbf{r}_2 = \eta_0 [\nabla \mathbf{v}_2 + (\nabla \mathbf{v}_2)^T] + \left[ \eta_0^b - \frac{\zeta}{2} \eta_0 \right] (\nabla \cdot \mathbf{v}_2) \mathbf{I}. \]  

(8b)

Using Eq. (1) in the form \( \rho_1 = \rho_0 \kappa_s \rho_1 \) and the first-order momentum equation (6a), we rewrite Eq. (8a) to

\[ \rho_0 \partial_t \mathbf{v}_2 = \nabla \cdot [\mathbf{r}_2 - \mathbf{p}_2 \mathbf{I}] - \kappa_s \mathbf{p}_1 \mathbf{r}_1 + \frac{1}{2} \kappa_s \mathbf{p}_1 \mathbf{r}_1 \mathbf{I} \]

\[ + \kappa_s \mathbf{v}_1 \cdot \mathbf{r}_1 - \rho_0 (\mathbf{v}_1 \cdot \nabla) \mathbf{v}_1. \]  

(8c)

This particular form of the second-order momentum equation is chosen to minimize numerical errors as described in Sec. III A.

E. Periodic frequency-domain equations

When solving for the periodic state at \( t \to \infty \), it is advantageous to formulate the first-order equations in the frequency domain. The harmonic first-order fields are all written as \( \mathbf{g}_1(r, t) = \text{Re}[\mathbf{g}_1^d(r)e^{-i\omega t}w] \), where \( \mathbf{g}_1^d \) is the complex field amplitude in the frequency domain. The first-order frequency-domain equations are derived from Eqs. (5) and (6a) by the substitution \( \partial_t \to -i \omega \).

\[ \nabla \cdot \mathbf{v}_1^d - i \omega \kappa_s \mathbf{p}_1^d = 0, \]  

(9)

\[ \nabla \cdot [\mathbf{r}_1^d - \mathbf{p}_1^d \mathbf{I}] + i \omega \rho_0 \mathbf{v}_1^d = \mathbf{0}. \]  

(10)

The steady state-averaged streaming flow is obtained from the time-averaged second-order frequency-domain equations, where \( \langle g^d \rangle \) denotes time averaging over one oscillation period of the periodic second-order field. The time average of products of two harmonic first-order fields \( \langle g_1^d g_2^d \rangle \) is given by \( \langle g_1^d g_2^d \rangle = \frac{1}{2} \text{Re} [\langle \mathbf{g}_1^d \rangle \langle \mathbf{g}_2^d \rangle^*] \), as in Ref. [11], where the asterisk denotes complex conjugation. In the periodic state, the fields may consist of harmonic terms and a steady term, and thus all full time derivatives average to zero \( \langle \partial_t \mathbf{g}_1^d \rangle = 0 \). The time-averaged second-order frequency-domain equations are derived from Eqs. (7) and (4a),

\[ \nabla \cdot [\mathbf{v}_2^d] + \kappa_s \mathbf{v}_1^d \cdot \nabla \mathbf{p}_1^d = \mathbf{0}, \]  

(11)

\[ \nabla \cdot [\langle \mathbf{r}_2^d \rangle - \langle \mathbf{p}_2^d \rangle \mathbf{I}] - \rho_0 \langle \mathbf{v}_1^d \mathbf{v}_1^d \rangle = \mathbf{0}. \]  

(12)

F. Acoustic energy and cavity Q factor

The total acoustic energy of the system in the time domain \( E_{ac}(t) \) and in the frequency domain \( E_{ac}^f(\omega) \) is given by

\[ E_{ac}(t) = \int_V \left[ \frac{1}{2} \kappa_s \rho_1^2 + \frac{1}{2} \rho_0 v_1^2 \right] dV. \]  

(13a)

\[ E_{ac}^f(\omega) = \int_V \left[ \frac{1}{2} \kappa_s \rho_1^1 \rho_1^1 + \frac{1}{2} \rho_0 \rho_0 \mathbf{v}_1^1 \mathbf{v}_1^1 \right] dV. \]  

(13b)

Moreover, the time derivative of \( E_{ac}(t) \) is

\[ \partial_t E_{ac} = \int_V \left[ \kappa_s \rho_1 \partial_t \rho_1 + \rho_0 \mathbf{v}_1 \cdot \partial_t \mathbf{v}_1 \right] dV \]

\[ = \int_V [\mathbf{v}_1 \cdot (\mathbf{r}_1 - \mathbf{p}_1 \mathbf{I})] - \mathbf{v}_1 \cdot \partial_t \mathbf{v}_1 ] dV. \]  

(14a)

where we have used Eqs. (5) and (6a). Applying Gauss’s theorem on the first term in Eq. (14a), we arrive at

\[ \partial_t E_{ac} = \int_A [\mathbf{v}_1 \cdot (\mathbf{r}_1 - \mathbf{p}_1 \mathbf{I})] \cdot \mathbf{n} dA - \int_V \nabla \mathbf{v}_1 : \partial_t \mathbf{v}_1 dV \]

\[ = P_{\text{pump}} - P_{\text{dis}}, \]  

(14b)

where \( P_{\text{pump}} \) is the total power delivered by the forced vibration of the sidewalls, and \( P_{\text{dis}} \) is the total power dissipated due to the viscous stress. The quality factor \( Q \) of a resonant cavity is given by

\[ Q = \frac{2\pi E_{ac}^f}{\text{Energy dissipated per cycle}} = \frac{1}{\omega} \left( \frac{E_{ac}^f}{P_{\text{dis}}} \right). \]  

(15)

G. Summary of theory

Throughout this paper we refer to two kinds of solutions of the acoustic energy and velocity fields: unsteady nonperiodic solutions obtained from Eqs. (5)–(8) and steady periodic solutions obtained from Eqs. (9)–(12). When presenting the unsteady nonperiodic solutions, they are often normalized by the steady periodic solution, to emphasize how close it has converged towards this solution.

III. NUMERICAL MODEL

As summarized in Sec. I, the vast majority of experimental work in the field of acoustophoresis deals with frequencies around \( f \approx 10 \text{ MHz} \), acoustic pressures of the order \( p_1 \approx 0.1 \text{ MPa} \), boundary layer widths \( \delta < 500 \text{ nm} \), and channel cross sections of dimensions near \( h \approx 0.2 \text{ mm} \). The key dimensional numbers are then low Mach numbers \( \text{MA} = v_1/c_0 = p_1/(\rho_0 \omega_0^2) < 10^{-3} \), fairly low acoustic Reynolds numbers...
The vibrations in the surrounding silicon-glass chip and piezo transducer are not modeled. The water domain is surrounded by immovable hard walls, and the acoustic fields are excited on the side walls by an oscillating velocity boundary condition with oscillation period $t_0$ and angular frequency $\omega$,

$$t_0 = \frac{2\pi}{\omega}. \quad (16)$$

The symmetry of the bottom boundary is described by zero orthogonal velocity component and zero orthogonal gradient of the parallel velocity component. The explicit boundary conditions for the first-order velocity become

- top: $v_y = 0, \quad v_z = 0, \quad (17a)$
- bottom: $\partial_z v_y = 0, \quad v_z = 0, \quad (17b)$
- left-right: $v_y = v_{bc} \sin(\omega t), \quad v_z = 0. \quad (17c)$

The boundary conditions on the second-order velocity are set by the zero-mass-flux condition $\mathbf{n} \cdot \mathbf{\rho} = 0$ on all boundaries, as well as zero parallel velocity component on the top, right, and left wall boundaries, and zero orthogonal derivative of the parallel component of the mass flux on the bottom symmetry boundary. The explicit boundary conditions for the second-order velocity become

- top: $v_{yz} = 0, \quad v_z = 0, \quad (18a)$
- bottom: $\partial_z (\mathbf{\rho}_0 v_{yz} + \mathbf{\rho}_l v_y) = 0, \quad v_z = 0, \quad (18b)$
- left-right: $\mathbf{\rho}_0 v_{yz} + \mathbf{\rho}_l v_y = 0, \quad v_z = 0. \quad (18c)$
C. Spatial resolution

The physical fields are discretized using fourth-order basis functions for \( v_x \) and \( v_y \) and third-order basis functions for \( p_x \) and \( p_z \). The domain shown in Fig. 1(a) is covered by basis functions localized in each element of the spatial mesh shown in Fig. 1(c). Since the streaming flow is solved in the time domain, the computational time quickly becomes very long compared to the computational time of solving the usual steady streaming flow. Thus we have optimized the use of precious few mesh elements to obtain the best accuracy of the solution. We use an inhomogeneous mesh of rectangular elements ranging in size from 0.16 \( \mu \text{m} \) at the boundaries to 24 \( \mu \text{m} \) in the bulk of the domain. The convergence of the solution \( g \) with respect to a reference solution \( g_{\text{ref}} \) was considered through the relative convergence parameter \( C(g) \) defined in Ref. [11] by

\[
C(g) = \sqrt{\frac{\int (g - g_{\text{ref}})^2 dy dz}{\int (g_{\text{ref}})^2 dy dz}}. \tag{19}
\]

In Ref. [11], \( C(g) \) was required to be below 0.001 for the solution to have converged. The solution for the steady time-averaged velocity \( \langle v_x^2 \rangle \), calculated with the mesh shown in Figs. 1(c) and 1(d), has \( C = 0.006 \) with respect to the solution calculated with the fine triangular reference mesh in Ref. [11], which is acceptable for the present study.

D. Temporal resolution

The required temporal resolution for time-marching schemes is normally determined by the Courant-Friedrichs-Lewy (CFL) condition [31], also referred to as just the Courant number,

\[
\text{CFL} = \frac{c}{\Delta t} \leq \text{CFL}_{\text{max}}, \tag{20}
\]

where \( \Delta t \) is the temporal discretization and \( \Delta r \) is the spatial discretization. This means that the length over which a disturbance travels within a time step \( \Delta t \) should be some fraction of the mesh element size, ultimately ensuring that disturbances do not travel through a mesh element in one time step. A more accurate interpretation of the CFL condition is that it ensures that the error on the approximation of the time derivative is smaller than the error on the approximation of the spatial derivatives. Consequently, the value of CFL_{max} depends on the specific solver and on the order of the basis functions. For fourth-order basis functions and the generalized alpha solver, Ref. [31] reports a value of CFL_{max} = 0.05, which is an empirical result for a specific model. Due to the inhomogeneity of the mesh, two values for the upper limit for the temporal resolution can be calculated based on Eq. (20): \( \Delta t = 8 \times 10^{-10} \text{ s} \approx 100/600 \) for the bulk mesh size of 24 \( \mu \text{m} \) and \( \Delta t = 5 \times 10^{-12} \text{ s} \approx 10^5/95 000 \) for the boundary mesh size of 160 nm.

To determine a reasonable trade-off between numerical accuracy and computational time, we study the convergence of the transient solution towards the steady solution for different values of the temporal resolution \( t_0/\Delta t \). The acoustic energy \( E_{\text{ac}}(t) \) is shown in Fig. 2(a) for different values of \( \Delta t \) and normalized by the steady time-averaged energy \( \langle E_{\text{ac}}^\infty (\infty) \rangle \) of the frequency-domain calculation, and it is thus expected to converge to unity for long times. In Fig. 2(b), \( E_{\text{ac}}(1000 t_0)/\langle E_{\text{ac}}^\infty (\infty) \rangle \) is plotted versus the temporal resolution \( t_0/\Delta t \), which shows how the accuracy of the time-domain solution increases as the temporal resolution is increased. In all subsequent simulations we have chosen a time step of \( \Delta t = t_0/256 \), the circled point in Fig. 2(b), for which the time-domain energy converges to 99.4% of the energy of the steady calculation. The chosen value for the time step is larger than the upper estimate \( t_0/600 \) of the necessary \( \Delta t \) based on the CFL condition. This might be because our spatial domain is smaller than the wavelength, and consequently a finer spatial resolution is needed, compared to what is usually expected to spatially resolve a wave.

We have noted that the fastest convergence is obtained when actuating the system at its (numerically determined) resonance frequency \( f_{\text{res}} \). When shifting the actuation frequency half the resonance width \( 1/2 \Delta f \) away from \( f_{\text{res}} \), the energy \( E_{\text{ac}}(t) \) for \( \Delta t = t_0/256 \) converged to only 95% of the steady value \( \langle E_{\text{ac}}^\infty (\infty) \rangle \) (calculated in the frequency domain), thus necessitating smaller time steps to obtain reasonable convergence.

The computations where performed on a desktop PC with Intel Xeon CPU X5690 3.47 GHz 2 processors, 64-bit...
IV. ONSET OF ACOUSTIC STREAMING

In this section the fluid is initially quiescent. Then, at time \( t = 0 \), the oscillatory velocity actuation is turned on, such that within the first oscillation period its amplitude increases smoothly from zero to its maximum value \( v_{bc} \), which it maintains for the rest of the simulation. We study the resulting buildup of the acoustic resonance and the acoustic streaming flow.

A. Resonance and buildup of acoustic energy

To determine the resonance frequency, the steady acoustic energy \( \langle E_{\text{ac}}(\infty) \rangle \) Eq. (13b) was calculated for a range of frequencies based on the frequency-domain equations (9) and (10). In Fig. 3 the numerical results (circles) are shown together with a Gaussian fit (full line), while the inset exhibits the fitted resonance frequency \( f_{\text{res}} \), the full width \( \Delta f \) at half maximum, and the quality factor \( Q = f_{\text{res}}/\Delta f \).

The buildup of the acoustic energy in the cavity is well captured by a simple analytical model of a single sinusoidally driven damped harmonic oscillator with time-dependent position \( x(t) \),

\[
\frac{d^2x}{dt^2} + 2\Gamma_0 x + \omega_0^2 x = \frac{1}{m} F_0 \sin(\omega t), \tag{21}
\]

\[
\frac{d^2x}{dt^2} + 2\Gamma_0 x + \omega_0^2 x = \frac{1}{m} F_0 \sin(\omega t).
\]

The amplitude \( A \) and the phase shift \( \phi \) between the forcing and the response are known functions of \( \frac{\omega}{\omega_0} \), \( \omega_0 \), \( \omega \), and \( \Gamma \), which are not relevant for the present study. From Eq. (22) we obtain the velocity \( dx/dt \), leading to the total energy \( E \) of the oscillator,

\[
\frac{1}{2} m \left[ \frac{d}{dt} x(t) \right]^2 + \frac{1}{2} Q v \left[ \frac{d}{dt} x(t) \right]^2. \tag{23}
\]

The amplitude \( A \) and the phase shift \( \phi \) between the forcing and the response are known functions of \( \frac{\omega}{\omega_0} \), \( \omega_0 \), \( \omega \), and \( \Gamma \), which are not relevant for the present study. From Eq. (22) we obtain the velocity \( dx/dt \), leading to the total energy \( E \) of the oscillator,

\[
E(t) = \frac{1}{2} m \omega_0^2 A^2 + \frac{1}{2} Q \left[ \frac{d}{dt} A(t) \right]^2.
\]

Based on Eqs. (22) and (23), the characteristic time scale \( \tau_E \) for the buildup of the acoustic energy is found to be

\[
\tau_E = \frac{1}{2 \Gamma_0} = \frac{Q}{\omega_0}. \tag{24}
\]

The buildup of the energy in the single harmonic oscillator, calculated at \( \omega = \omega_0 \) with \( \Gamma = 1.20 \times 10^{-3} \), is shown in the inset of Fig. 3 together with the buildup of acoustic energy \( E(t) \) of the microfluidic channel solved numerically at resonance, \( \omega = 2\pi f_{\text{res}} \). The analytical and numerical results are in good agreement, and we conclude that the buildup of acoustic energy in the channel cavity can be modeled as a single harmonic oscillator. The energy builds up to 95% of its steady value in about 500 \( t_0 \) \( \approx 8 \tau_E \).

B. Decomposition of the velocity field

The task of calculating the buildup of the acoustic streaming flow is a multiscale problem, because the amplitude of the oscillating acoustic velocity field is several orders of magnitude larger than the magnitude of the streaming flow. This is indeed the very reason that we can apply the perturbation expansion

\[
v = v_1 + v_2, \tag{25}
\]

and decompose the nonlinear governing equations into a set of linear first-order equations and a set of second-order equations. However, there is also another level of difference in velocity scaling. In the purely periodic state, the velocity can be Fourier decomposed as

\[
v(r,t) = v_1(r) \sin(\omega t) + v_2(r) \sin(2\omega t) + v_0(r), \tag{26}
\]

where \( v_1(r) \) is the steady amplitude of the first-order harmonic component, \( v_2(r) \) is the steady amplitude of the second-order frequency-doubled component, and \( v_0(r) \) is the magnitude of the second-order steady velocity component referred to as the acoustic streaming flow. The orders of magnitude of the three velocity components in the periodic state are given by

\[
v_1 \approx Q v_{bc}, ~ v_2 \approx Q^2 v_{bc}, ~ v_0 \approx Q^2 v_{bc}. \tag{27}
\]
The order of \( v_1 \) is derived in the one-dimensional acoustic cavity example presented in Ref. [32], the order of \( v_2^0 \) is given by the well-known Rayleigh theory, while the order of \( v_2^{20} \) is derived in the Appendix. The magnitude of \( v_2^{20} \) is a factor of \( Q \) larger than what is expected from dimensional analysis of the second-order equation (8c). Consequently, the criterion \( |v_2| \ll |v_1| \) for the perturbation expansion becomes

\[
Q^2 v_{bc} \ll c_s, \tag{28}
\]

which is more restrictive than the usual criterion based on the first-order perturbation expansion, \( Q v_{bc} \ll c_s \). Thus, the perturbation expansion becomes invalid for smaller values of \( v_{bc} \) than previously expected.

In the transient regime we cannot Fourier decompose the velocity field. Instead, we propose a decomposition using envelope functions inspired by Eq. (26),

\[
\mathbf{v}(\mathbf{r}, t) = v_0^0(r, t) \sin(\omega t) + \mathbf{v}_2^0(r, t) \sin(2\omega t) + v_2^0(r, t).
\]

Here, the amplitudes are slowly varying in time compared to the fast oscillation period \( t_0 \). We can no longer separate \( v_2^{20} \) and \( v_0^0 \) before solving the second-order time-dependent equations (7) and (8). To obtain the time-dependent magnitude of the quasisteady streaming velocity mode \( v_0^0 \), we need to choose a good velocity probe, and we thus form the unsteady time average of \( v_2(r, t) \),

\[
\langle v_2(r,t) \rangle = \int_{t-t_0/2}^{t+t_0/2} v_2(r,t')dt'. \tag{30}
\]

The time averaging is done with a fifth-order Romberg integration scheme [33] using data points with a uniform spacing of \( t_0/16 \) in the time interval of width \( t_0 \).

**C. Steady and unsteady streaming flow**

In this section we compare the unsteady time-averaged second-order velocity field \( \langle \mathbf{v}_2(r,t) \rangle \), from the time-domain simulations, with the steady time-averaged second-order velocity field \( \langle \mathbf{v}_2^0(\infty) \rangle \), from the frequency-domain simulation. Figures 4(a) and 4(b) each show a snapshot in time of the transient \( \mathbf{v}_1 \) and \( \mathbf{v}_2 \), respectively. For \( \mathbf{v}_2(r,t) \), the oscillatory component \( v_2^{20}(r,t) \sin(2\omega t) \) dominates, as it is two orders of magnitude larger than the quasisteady component \( v_2^0(r,t) \). However, at late times, here \( t = 3000 t_0 \), the amplitude \( \mathbf{v}_2^{20}(r,t) \) has converged, and in \( \langle \mathbf{v}_2(r,t) \rangle \) the oscillatory component averages to zero and only the quasisteady component remains.

The unsteady time average \( \langle v_2(r,t) \rangle \) evaluated at \( t = 3000 t_0 \) is shown in Fig. 4(c), exhibiting a single flow roll, in agreement with the classical Rayleigh streaming flow.

In Fig. 4(d) is shown the steady \( \langle \mathbf{v}_2^0(\infty) \rangle \) from the frequency-domain simulation. Figures 4(c) and 4(d) use the same color scaling for the velocity magnitude, to evaluate the convergence of the unsteady streaming flow \( \langle \mathbf{v}_2(3000 t_0) \rangle \) towards the steady streaming flow \( \langle \mathbf{v}_2^0(\infty) \rangle \), and the two solutions agree well both qualitatively and quantitatively.

The convergence parameter \( C \), Eq. (19), of \( \langle v_2(3000 t_0) \rangle \) with respect to \( \langle \mathbf{v}_2^0(\infty) \rangle \) is \( C = 0.01 \), and if we multiply \( \langle v_1 \rangle \) by a free factor, taking into account that \( \langle v_2 \rangle \) has not fully converged at \( t = 3000 t_0 \), the convergence parameter can be reduced to \( C = 0.008 \). The remaining small difference between the unsteady \( \langle v_2(3000 t_0) \rangle \) and the steady \( \langle \mathbf{v}_2^0(\infty) \rangle \) is attributed to the finite temporal resolution of the time marching scheme. We can thus conclude that the time-domain streaming simulation converges well towards the frequency-domain simulation, and this constitutes the primary validation of the unsteady nonperiodic simulations.
D. Buildup of the velocity field

To visualize the buildup of the acoustic fields over short and long time scales, we have chosen the three point probes shown in Fig. 1(b). The oscillating first-order velocity field is probed in the center of the channel \((0,0)\), far from the walls in order to measure the bulk amplitude of the acoustic field. The horizontal component of the second-order velocity \(v_2^y\) is probed on the horizontal symmetry axis at \((\frac{1}{4}w,0)\), where the oscillatory component \(v_2^y\omega^2\) has its maximum amplitude. The vertical component of the second-order velocity \(v_2^z\) is probed on the vertical symmetry axis at \((0,\frac{1}{4}h)\) where the oscillatory component \(v_2^z\omega^2\) is small and of the same order as the quasisteady component \(v_2^0\), making the unsteady time-averaged second-order velocity at this point a good probe for the quasisteady streaming velocity.

In Fig. 5 is shown the buildup of the velocity probes [panels (a)–(c)] and their time averages [panels (d)–(f)] for the first 20 oscillations. The thick lines are the oscillating velocities, while the thin lines are the envelopes of the oscillations. Already within the first 20 oscillation periods we see in Fig. 5(f) the buildup of a quasisteady velocity component. The time average of \(v_2^y\) tends to zero for long times as it is purely oscillatory, whereas the time average of \(v_2^z\) tends to the magnitude of the quasisteady component \(v_2^0\), because the large but now steady oscillatory component \(v_2^0\omega^2\) averages to zero. The dashed lines in Figs. 6(e) and 6(f) represent the magnitude of the steady time-averaged second-order velocity \(\langle v_2^z(\infty) \rangle\) from the frequency-domain simulation.

V. ACOUSTIC STREAMING GENERATED BY PULSED ACTUATION

In the following we study the effects of switching the oscillatory boundary actuation on and off on a time scale much longer than the oscillation period \(t_0\) in either single- or multipulse mode. The aim is to investigate whether such an approach can suppress the influence of the streaming flow on suspended particles relative to that of the radiation force.

A. Single-pulse scaling analysis

A striking feature of Fig. 6 is the separation of time scales between the roughly exponential buildup of the acoustic resonance in Fig. 6(a) and of the streaming flow in Fig. 6(f). It appears that the resonance, and hence the acoustic radiation force on a suspended particle, is fully established almost ten times faster than the streaming flow and the resulting drag force on a suspended particle. To investigate this further, we look at the scaling provided by the three time scales relevant for the problem of transient acoustic streaming, all listed in Table II: the oscillation time \(t_0\) of the acoustic wave, the resonance magnitude \(v_1^0/c_s \sim 3 \times 10^{-4}\) [Fig. 6(a)], \(v_2^0/c_s \sim 5 \times 10^{-5}\) [Fig. 6(b)], and \(v_2^0/c_s \sim 1 \times 10^{-7}\) [Figs. 6(e) and 6(f)]. The time average of \(v_2^y\) tends to zero for long times as it is purely oscillatory, whereas the time average of \(v_2^z\) tends to the magnitude of the quasisteady component \(v_2^0\), because the large but now steady oscillatory component \(v_2^0\omega^2\) averages to zero. The dashed lines in Figs. 6(e) and 6(f) represent the magnitude of the steady time-averaged second-order velocity \(\langle v_2^z(\infty) \rangle\) from the frequency-domain simulation.
relaxation time $\tau_E$ of the acoustic cavity, and the momentum diffusion time $\tau_v$ governing the quasisteady streaming flow.

The momentum diffusion time is $\tau_v = \frac{1}{2\nu}(\frac{h}{\nu})^2$, where $\nu = \frac{\eta}{\rho_0}$ is the kinematic viscosity, and $\frac{1}{2}h$ is approximately half the distance between the top boundary and the center of the streaming flow roll. Inserting the relevant numbers, see Table II, we indeed find that $\tau_E \approx 66t_0$ is much faster than $\tau_v \approx 558t_0$.

While these numbers are obtained for our specific choice of parameters, we can easily establish the scaling for the ratio $\tau_v/\tau_E$ as a function of the system parameters in general. For the momentum diffusion time we have $\tau_v \propto h^2/\nu$, while the resonance relaxation time at the half-wavelength resonance is $\tau_E \propto Q(t_0) \propto \frac{1}{\lambda^2} \propto \frac{h}{\sqrt{\nu}} \propto h/\sqrt{\nu}$. Thus $\tau_v/\tau_E \propto \frac{h}{\sqrt{\nu} \nu}$. and the separation of time scales can be increased in this case by increasing the channel height $h$, decreasing the channel thickness $w$, and decreasing the kinematic viscosity $\nu$.

However, this separation in time scales does not guarantee a suppression of streaming relative to the radiation force. One problem is that the streaming is shear-stress driven in the momentum diffusion time $\tau_v$, and not dependent on the small thickness of the boundary layer. This we investigate further in the following subsection. Another problem is that the large momentum diffusion time $\tau_v$ implies a very slow decay of the streaming flow, once it is established. The latter effect, we study using the following analytical model. Consider a quantity $f$ (streaming velocity or acoustic energy), with a relaxation time $\tau$ and driven by a pulsed source term $P$ of pulse width $t_{pw}$. The rate of change of $f$ is equivalent to Eq. (14b),

$$\frac{\partial f}{\partial t} = P - \frac{1}{\tau} f,$$

(31a)

$$P = \begin{cases} \frac{1}{\tau} f_0, & \text{for } 0 < t < t_{pw}, \\ 0, & \text{otherwise,} \end{cases}$$

(31b)

where $\frac{1}{\tau} f_0$ is a constant input power. This simplified analytical model captures the roughly exponential buildup and decay characteristics of our full numerical model, and allows for analytical studies of the time integral of $f(t)$. For a finite time $t > t_{pw}$, we find

$$\int_0^t f(t') dt' = f(t_{pw}) - f_0 \left[ e^{-\frac{t}{\tau}} - e^{-\frac{t_{pw}}{\tau}} \right].$$

(32)

From this we see that when $t \gg \tau + t_{pw}$ the time-integral of $f(t)$ is approximately $f(t_{pw})$ and not dependent on the relaxation time $\tau$. Consequently, if both the acoustic energy and the acoustic streaming can be described by exponential behavior with the respective relaxation times $\tau_E$ and $\tau_v$, the ratio of their time-integrated effects is the same whether the system is driven by a constant actuation towards their steady time-periodic state or by a pulsed actuation with pulse width $t_{pw}$. This simplified analytical model indicates that there is little
and the unsteady energy and streaming probes obtained from the time-domain simulation are normalized by their corresponding steady time-averaged values from the frequency-domain simulation.

We introduce the streaming ratio \( \chi \) to measure the influence of streaming-induced drag on suspended particles relative to the influence of the acoustic radiation force for the unsteady time-domain solution, in comparison to the periodic frequency-domain solution. To calculate the relative displacement \( \Delta s \) of particles due to each of the two forces, respectively, we compare their time integrals. Since the radiation force scales with the acoustic energy density, we define the streaming ratio \( \chi(t) \) as

\[
\chi(t) = \frac{\int_0^t \frac{\langle v_{\text{str}}(t') \rangle}{\langle E_{\text{ac}}(\infty) \rangle} dt'}{\int_0^t \frac{\langle v_{\text{rad}}(t') \rangle}{\langle E_{\text{ac}}(\infty) \rangle} dt'} \sim \frac{\Delta s_{\text{str}}}{\Delta s_{\text{rad}}},
\]

where \( \Delta s_{\text{str}} \) and \( \Delta s_{\text{rad}} \) are the total particle displacements in the time from 0 to \( t \) due to the streaming-induced drag force and the acoustic radiation force, respectively. In the periodic state \( \chi = 1 \), and to obtain radiation-force-dominated motion of smaller particles, we need to achieve a smaller value of \( \chi \). Obtaining a value of \( \chi = 0.8 \) at time \( t_{\text{end}} \) implies that the ratio of the relative displacement due to the streaming-induced drag force and the radiation force for the time interval \( 0 < t < t_{\text{end}} \) is 20% lower than in the periodic state, corresponding to a 20% reduction of the critical particle size for acoustophoretic focusing, defined in Ref. [2], assuming the particles can be focused during the time interval \( 0 < t < t_{\text{end}} \).

Figure 7(a) shows \( \langle E_{\text{ac}} \rangle \), \( \langle v_{\text{str}} \rangle \), and \( \chi \) during the buildup towards the periodic state. \( \chi \) approaches unity slower than \( \langle v_{\text{str}} \rangle \) because \( \chi \) is an integration of the streaming and radiation contributions, whereas \( \langle v_{\text{str}} \rangle \) probes the instantaneous magnitude of the streaming flow. Figure 7(b) and 7(c) show \( \langle E_{\text{ac}} \rangle \), \( \langle v_{\text{str}} \rangle \), and \( \chi \) when the actuation is turned off at \( t = 200 t_0 \) and \( t = 30 t_0 \), respectively. When the actuation is turned off, \( \langle E_{\text{ac}} \rangle \) decays faster than \( \langle v_{\text{str}} \rangle \) and thus \( \chi \) begins to increase more rapidly, reaching \( \chi = 0.8 \) around \( t = 1000 t_0 \) in both cases. From the results shown in Fig. 7 it does not seem advantageous to turn off the actuation, as this only causes \( \chi \) to increase faster than for constant actuation. Figure 7(c) further shows that when the actuation is turned off, \( \langle E_{\text{ac}} \rangle \) immediately begins to decay, whereas \( \langle v_{\text{str}} \rangle \) continues to increase for some time, due to the present acoustic energy in the system that still provides a driving force for the streaming flow.

B. Single-pulse numerical simulation

We investigate the features of pulsed actuation in more detail in the following by numerical simulation. In Fig. 7 is shown the temporal evolution of the total acoustic energy \( \langle E_{\text{ac}} \rangle \) and the magnitude of the acoustic streaming flow \( \langle v_{\text{str}} \rangle \) for the three cases: (i) the buildup towards the periodic state, (ii) a single long actuation pulse, and (iii) a single short actuation pulse. The magnitude of the acoustic streaming is measured by the unsteady time-averaged velocity probe

\[
\langle v_{\text{str}} \rangle = \langle v_z(0, \frac{1}{2}h) \rangle,
\]

and

\[
\langle v_{\text{str}} \rangle = \frac{1}{t_{\text{end}}} \int_0^{t_{\text{end}}} \langle v_z(t) \rangle dt,
\]

where \( t_{\text{end}} \) is the duration of the actuation pulse.

C. Multipulse numerical simulation

From the single-pulse results shown in Fig. 7 there is no indication of any optimum for the pulse duration or repetition period, and in general it provides little hope that pulsed actuation should lead to lower values of \( \chi \). Figure 8 shows \( \langle E_{\text{ac}} \rangle \), \( \langle v_{\text{str}} \rangle \), and \( \chi \) for three pulsed schemes with pulse duration 500 \( t_0 \), 200 \( t_0 \), and 30 \( t_0 \) and pause duration 500 \( t_0 \), 200 \( t_0 \), and 210 \( t_0 \), respectively. For all three pulsed schemes, \( \chi \) increases faster than for the constant actuation Fig. 7(a), thus not indicating any increased suppression of the streaming.
numerical time-domain solver, which in the generalized-alpha solver is done through the alpha parameter. Simulation with higher temporal resolution required lower values of the alpha parameter to have more suppression of accumulated numerical errors.

The model system used in this study is a simplification of an actual device. The vibration of only the side walls, and not the top and bottom walls, stands in contrast to the physical system, in which the whole device is vibrating in a nontrivial way, difficult to predict, and only the overall amplitude and the frequency of the actuation is controlled experimentally. Furthermore, our model only treats the two-dimensional cross section of a long straight channel, whereas experimental studies have shown that there are dynamics along the length of the channel [21]. Nevertheless, successful comparisons, both qualitatively and quantitatively, have been reported between the prediction of this simplified numerical model and experimental measurements of Rayleigh streaming in the cross-sectional plane of a microchannel [10], which makes it reasonable to assume that the time-dependent simulations also provide reliable predictions.

It is also important to stress that our model only describes the fluid and not the motion of the suspended particles. Integrating the forces acting on the particles becomes vastly more demanding when the streaming flow is unsteady, because the drag forces from the oscillating velocity components $v_1$ and $v_2 \omega$ do not average out, as they do in the case of a purely time-periodic state. To include this contribution in the particle tracking scheme, the forces on the particles need to be integrated with a time step of a fraction of the oscillation period, which makes the solution of particle trajectories over several seconds a very demanding task using brute-force integration of the equations of motion.

Our analysis of the pulsed actuation schemes showed that the slow decay of the streaming flow makes pulsation inefficient in reducing the streaming-induced drag force compared to the radiation force. Such a reduction may, however, be obtained by a rapid switching between different resonances each resulting in similar radiation forces but different spatial streaming patterns which on average cancel each other out, thus fighting streaming with streaming. An idea along these lines was presented by Ohlin et al. [34], who used frequency sweeping to diminish the streaming flows in liquid-filled wells in a multiwell plate for cell analysis. However, the prediction of particle trajectories under such multiresonance conditions requires an extensive study as described above.

To our knowledge, experimental studies of pulsed actuation to decrease streaming flow have so far only been reported in the literature by Hoyos et al. [18]. Unfortunately, their study is not directly comparable to our analysis, as we treat the buildup of Rayleigh streaming perpendicular to the pressure nodal plane, whereas Hoyos et al. studied the streaming flow in this plane. Such in-nodal-plane streaming flows have been studied numerically by Lei et al. [12,13], though only with steady actuation. The contradicting results of our numerical study and the experimental study of Hoyos et al. may thus rely on the differences of the phenomena studied. However, despite the negative result presented above, we nevertheless hope that our analysis may serve as an inspiration for future experimental and numerical studies of acoustic streaming induced by pulsed

VI. DISCUSSION

Solving numerically the time-dependent problem of the acoustic cavity and the buildup of acoustic streaming presents new challenges which are not present in the purely periodic problem. First, the numerical convergence analysis now involves both the spatial and temporal resolutions. This we addressed in a sequential process by first analyzing the spatial mesh with the periodic frequency-domain solution, and thereafter doing a thorough convergence analysis with respect to the temporal resolution. Second, the convergence of the transient solution towards the periodic state was poor for actuation frequencies away from the resonance frequency of the system. This makes off-resonance simulation computationally costly, as it requires a better temporal resolution, and it complicates comparison of simulations at resonance with simulations off resonance. Third, small numerical errors accumulate during the hundred thousand time steps taken during a simulation from a quiescent state to a purely periodic state. These errors need to be suppressed by the

FIG. 8. (Color online) The same probes as in Fig. 7 but for the following pulsed actuation schemes: (a) actuation is on for $500 \tau_0$ followed by no actuation for $500 \tau_0$ repeatedly, (b) actuation is on for $200 \tau_0$ followed by no actuation for $200 \tau_0$ repeatedly, and (c) actuation is on for $30 \tau_0$ followed by no actuation for $210 \tau_0$ repeatedly.
actuation. We would certainly be interested in applying our analysis to future, well-characterized experimental studies of such pulse-actuated streaming.

VII. CONCLUSION

In this work, we have presented a model for the transient acoustic fields and the unsteady time-averaged second-order velocity field in the transverse cross-sectional plane of a long straight microchannel. The model is based on the usual perturbation approach for low acoustic field amplitudes, and we have solved both first- and second-order equations in the time domain for the unsteady transient case as well as in the frequency domain for the purely periodic case. This enabled us to characterize the buildup of the oscillating acoustic fields and the unsteady streaming flow.

Our analysis has shown that the buildup of acoustic energy in the channel follows the analytical prediction obtained for a single damped harmonic oscillator with sinusoidal forcing, and that a quasisteady velocity component is established already within the first few oscillations and increases in magnitude as the acoustic energy builds up. We have also found that for a resonance with quality factor $Q$, the amplitude of the oscillatory second-order velocity component is a factor of $Q$ larger than what is expected from dimensional analysis, which results in a more restrictive criterion for the validity of the perturbation expansion, compared to the usual one based on the first-order perturbation expansion.

Furthermore, contrary to a simple scaling analysis of the time scales involved in the fast buildup of radiation forces and slow buildup of drag-induced streaming forces, we have found that pulsating oscillatory boundary actuation does not reduce the time-integrated streaming-induced drag force relative to the time-integrated radiation force. As a result, pulsating actuation does not prevent streaming flows perpendicular to the pressure nodal plane from destroying the ability to focus small particles by acoustophoresis.

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APPENDIX: AMPLITUDE OF THE SECOND-ORDER OSCILLATORY VELOCITY FIELD

Extending to second order the one-dimensional example given in Ref. [32], we derive in this Appendix the order of magnitude of the second-order oscillatory component $v_2^{2o}$, which was stated in Eq. (27).

Just as $(g_2)$ denotes time-averaging over one oscillation period, Eq. (30), and in the periodic state equals the zero-order temporal Fourier component of the field, $g_2^{2o}(r)$ denotes the complex amplitude of the oscillatory second-order mode and is given by the second-order Fourier component,

$$g_2^{2o}(r) = \frac{1}{T} \int_{-T/2}^{+T/2} g_2(r, t') e^{-i2\omega t'} dt'.$$  \hspace{1cm} (A1)

By using that for any complex number $Z$ we have $\text{Re}[Z] = \frac{1}{2}(Z + Z^*)$, the product $A(r, t)B(r, t)$ of two oscillating fields $A(r, t) = \text{Re}[A e^{-i\omega t}]$ and $B(r, t) = \text{Re}[B e^{-i\omega t}]$ can be decomposed into a steady component and an oscillatory component,

$$A(t)B(t) = \frac{1}{2}(A e^{-i\omega t} + A^* e^{i\omega t})(B e^{-i\omega t} + B^* e^{i\omega t})$$
$$= \frac{1}{2}\text{Re}[A^* B] + \frac{1}{2}\text{Re}[A B e^{-i2\omega t}].$$  \hspace{1cm} (A2)

from which we introduce the following notation:

$$\langle AB \rangle \equiv \frac{1}{2}\text{Re}[A^* B], \quad (AB)^{2o} \equiv \frac{1}{2}AB,$$  \hspace{1cm} (A3)

where $A$ and $B$ could be any first-order fields.

The governing equations for the oscillatory second-order component $v_2^{2o}$ can be derived from Eqs. (7) and (8), and in the one-dimensional problem treated in Ref. [32], where the top and bottom walls are not taken into account, they become

$$-i2\omega \kappa_p v_2^{2o} = -\partial_y v_2^{2o} - \kappa_0 (v_1 \partial_y p_1)^{2o},$$  \hspace{1cm} (A4a)

$$-i2\omega \rho_0 v_2^{2o} = -\partial_y p_2^{2o} - (\partial_y^2 + \eta^2) v_2^{2o} - [\rho_1 (-i\omega v_1)]^{2o} - [\rho_0 (v_1 \partial_y v_1)]^{2o}.$$  \hspace{1cm} (A4b)

Applying the $2\omega$ rule $(A3)$ and mass continuity $(5)$, the two last terms of Eq. (A4b) cancel. Inserting Eq. (A4a) into Eq. (A4b), the governing equation for $v_2^{2o}$ becomes

$$4k_0^2 v_2^{2o} + (1 - 4\Gamma\partial_y) v_2^{2o} + \frac{1}{2}\kappa_p \partial_y (v_1 \partial_y p_1) = 0,$$  \hspace{1cm} (A5)

where $\Gamma$ is the nondimensional bulk damping coefficient given by $\Gamma = \frac{\eta}{2\kappa_p c_s^2} (\frac{1}{\kappa_p} + \frac{\rho_0}{\rho_1})$, and $k_0 = \frac{\omega}{c_s}$ is the wave number. For the fundamental half-wave resonance, the spatial dependence of the source term $\partial_y (v_1 \partial_y p_1)$ is $\sin(2k_0 y)$, and the guess for the inhomogeneous solution to Eq. (A5) thus becomes

$$v_2^{2o, inhom} = C \sin(2k_0 y).$$  \hspace{1cm} (A6)

Inserting the inhomogeneous solution Eq. (A6) into the governing equation (A5), we note that the first term cancels with the “1” in the parentheses of the second term, and the order of magnitude of the inhomogeneous solution thus becomes

$$|v_2^{2o}| \sim C \sim \frac{1}{\Gamma} \kappa_p |v_1||p_1| \sim \frac{1}{\Gamma^1} \frac{v_2^{2o}}{c_s^2} \sim Q \frac{v_2^{2o}}{c_s^2},$$  \hspace{1cm} (A7)

which is the result stated in Eq. (27).
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