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Continuum modeling of hydrodynamic particle–particle interactions in microfluidic high-concentration suspensions

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A continuum model is established for numerical studies of hydrodynamic particle–particle interactions in microfluidic high-concentration suspensions. A suspension of microparticles placed in a microfluidic channel and influenced by an external force, is described by a continuous particle-concentration field coupled to the continuity and Navier-Stokes equation for the solution. The hydrodynamic interactions are accounted for through the concentration dependence of the suspension viscosity, of the single-particle mobility, and of the momentum transfer from the particles to the suspension. The model is applied on a magnetophoretic and an acoustophoretic system, respectively, and based on the results, we illustrate three main points: (1) for relative particle-to-fluid volume fractions greater than 0.01, the hydrodynamic interaction effects become important through a decreased particle mobility and an increased suspension viscosity. (2) At these high particle concentrations, particle-induced flow rolls occur, which can lead to significant deviations of the advective particle transport relative to that of dilute suspensions. (3) Which interaction mechanism that dominates, depends on the specific flow geometry and the specific external force acting on the particles.

1 Introduction

The ability to sort and manipulate biological particles suspended in bio-fluids is a key element in contemporary lab-on-a-chip technology.1–7 Such suspensions often contain bioparticles in concentrations so high that the particle–particle interactions change the flow behavior significantly compared to that of dilute suspensions. The main goal of this paper is to present a continuum model of hydrodynamic interaction effects in particle transport of high-concentration suspensions moving through microchannels in the presence of external forces.

One important example of high-concentration suspensions is full blood in blood vessels with diameters in the cm to sub-millimeter range or in artificial microfluidic channels. State-of-the-art calculations of the detailed hydrodynamics of blood involve direct numerical simulations (DNS).8–11 DNS models can simultaneously resolve the deformation dynamics of the elastic cell membranes of up to several thousand red blood cells (RBCs) as well as the inter-cell hydrodynamics of the fluid plasma, all in full 3D.9 However, such calculations require a steep price in computational efforts, as large computer clusters running thousands of CPUs in parallel may be involved,9 and the computations may take days.8 The computational cost may be reduced significantly by continuum modeling. Lei et al.12 investigated the limit of small ratios $a/H$ of the particle radius $a$ and the container size $H$. They found that, to a good approximation, continuum descriptions are valid for $a/H \simlt 0.02$. This is used in the computationally less demanding approach called mixture theory, or theory of interacting continua, where blood flow is modeled as two superimposed continua, representing the plasma and RBCs.13,14 In a recent study, Kim et al.14 used the method to simulate blood flow in a microchannel, and the results compared favorably to experiments. However, it is debated how to correctly calculate the coupling between the two phases, as well as formulate the stress tensor for the solid phase RBCs. This is called the ‘closure problem’ of mixture theory.15

Another important example is the hydrodynamics of macroscopic suspensions of hard particles, such as in gravity-driven sedimentation. These systems have been described by effective-medium theories of the effective viscosity of the suspension and of the effective single-particle mobility. The textbook by Happel and Brenner16 summarizes work from 1920 to 1960, in particular unit-cell modeling and infinite-array modeling, both approaches imposing a regular distribution of particles. Later, Batchelor17 improved the effective-medium models by applying statistical mechanics for irregularly places particles. In 1982, Mazur and Van Saarloos18 developed the induced-force method for explicit
calculation of the hydrodynamic interactions between many rigid spheres. A few years later, this method was extended by Ladd\textsuperscript{19,20} to allow for a higher number of particles to be considered in an efficient manner. In 1990, Ladd\textsuperscript{21} used the induced-force method in high-concentration suspensions to calculate three central hydrodynamic transport coefficients: the effective suspension viscosity \( \eta \), the particle diffusivity \( D \), and the particle mobility \( \mu \). The obtained results compare well with both experiments and with preceding theoretical work.

While studies of high-density suspensions with active transport are very important in lab-on-a-chip systems, such as in electro-, magneto-, and acoustophoresis\textsuperscript{1,3,4,22} they have not been treated by the microfluidic DNS-models mentioned above. On the other hand, while the effective-medium theories do treat active transport, they have rarely been used to analyze microfluidic systems. In this work, we present a continuum model of high-density suspensions in microchannels exposed to an external force causing particle migration. Based on the results of Lei \textit{et al.}\textsuperscript{12} mentioned above, we are justified to use the computationally less demanding continuum description in not too small microfluidic channels, \( a/H \lesssim 0.02 \). Moreover, with the hydrodynamic transport coefficients obtained by Ladd,\textsuperscript{21} we can fairly easily write down the governing equations for high-density suspensions following the Nernst–Planck-like approach presented by Mikkelsen and Bruus,\textsuperscript{23} but here extended to include the hydrodynamic particle–particle interactions.

We establish a continuum effective-medium theory, in which the hydrodynamic interactions between particles in a high-density suspension flowing through a microchannel are modeled by including the concentration dependency of the effective suspension viscosity, the particle diffusivity and mobility, and the transfer of momentum from the particles to the suspension. We then illustrate the use of the theory by two examples: The first is strongly inspired by the magneto-phoretic device studied by Mikkelsen and Bruus.\textsuperscript{23} The second is taken from the emerging field of acoustofluidics\textsuperscript{24,25} and involves an acoustophoretic device, for which the dilute limit is well described,\textsuperscript{16–18} but where the behavior in the high-concentration limit is still poorly characterized.

### 2 Theory

We consider a microchannel at ambient temperature, such as the one sketched in Fig. 1. In a steady-state pressure-driven flow of average inlet velocity \( v_0 \), an aqueous suspension of microparticles, with the homogeneous concentration \( c_0 \) (number of particles per volume) at the inlet, is flowing through the microchannel. Only a section of length \( L \) is shown. The effects of side-walls, present in actual systems, are negligible for channels with a large width-to-height ratio. Localized in the channel is an active area, where the microparticles are subjected to a transverse single-particle migration force \( F_{\text{mig}} \). This force changes the particle distribution, resulting in an inhomogeneous distribution downstream towards the outlet. In this process, the particles are transported by advection, active migration, and diffusion, all of which are affected by hydrodynamic particle–particle interactions at high concentrations.

![Fig. 1. Sketch of the model system.](image)

**2.1 Model system**

As a simple generic model system, we consider a section of length \( L \) of a microchannel with constant rectangular cross-section of height \( H \) and width \( W \). In specific calculations we take \( L = 7H \) and study for simplicity the flat-channel limit \( W = 10H \), for which the system, in a large part of the center, is reasonably well approximated by a parallel-plate channel. We take the microparticles to be rigid spheres with diameter \( 2a \). Their local concentration (number per volume) is described by the continuous field \( c(r) \), but often it is practical to work with the dimensionless concentration (particle volume fraction),

\[
\phi(r) = \frac{4}{3} \pi a^3 c(r).
\] (1)

We are particularly interested in high inlet concentrations \( \phi_0 \) between 0.001 and 0.1, and for this range with the typical values \( 2a = 2 \mu \text{m} \) and \( H = 50 \mu \text{m} \), the number of particles in the volume \( L \times W \times H \) is between \( 10^3 \) and \( 10^6 \), which in practice renders direct numerical simulation of the suspension infeasible. We therefore employ continuum modeling in terms of the three continuous fields: the particle concentration \( c(r) \), the pressure \( p(r) \), and the velocity \( \nu(r) \) of the suspension.

**2.2 Forces and effective particle transport**

Our continuum model for the particle current density \( J \) includes diffusion (diff), advection (adv), and migration (mig),\textsuperscript{29} for which we include hydrodynamic particle–particle interactions through \( c(r) \),

\[
J = J_{\text{diff}} + J_{\text{adv}} + J_{\text{mig}} \tag{2a}
\]

\[
= -D(c)\nabla c + cv + c\mu(c)F_{\text{mig}} \tag{2b}
\]

\[
= -f_D(\phi)D_0 \nabla c + cv + c f_p(\phi)\mu_0 F_{\text{mig}}. \tag{2c}
\]
Here,

\[ D(c) = \chi_D(c)D_0 \quad \text{and} \quad \mu(c) = \chi_\mu(c)\mu_0 \]  

(2d)
is the effective single-particle diffusivity and mobility, respectively. \( \chi_D(c) \) and \( \chi_\mu(c) \) concentration-dependent correction factors described in more detail in section 2.5, while \( \mu_0 = (6\pi\eta_0\text{d})^{-1} \) and \( D_0 = k_B T \eta_0 \) is the dilute-limit mobility and diffusivity, respectively, where \( k_B \) is the Boltzmann constant, \( T \) is the temperature, and \( \eta_0 \) is the dilute-limit viscosity of the suspension. We do not in this work consider any effects of van der Waals interactions or steric interactions.\(^2\)\(^9\) Due to conservation of particles, the particle current density obeys a steady-state continuity equation,

\[ 0 = \nabla \cdot J. \]  

(3)

2.3 Effective viscosity and flow of the suspension

The flow of the suspension is governed by the continuity equation for an incompressible suspension with the force density \( \sigma \), the stress tensor \( \sigma \) therefore becomes

\[ \rho \frac{\partial c}{\partial t} = \nabla \cdot (\rho \nabla \phi) + c F_{\text{mig}}, \]  

(5a)

\[ \sigma(\phi) = -1p + \chi_\eta(c)\phi^{-1}\eta \nabla v + (\nabla v)^T, \]  

(5b)

where 1 is the unit tensor. Likewise, the steady-state mass continuity equation for an incompressible suspension with the stress tensor \( \sigma \) becomes

\[ 0 = \nabla \cdot \left[ \chi_\rho(c)\rho_0 \nabla v \right]. \]  

(6)

2.4 Boundary conditions

We assume a concentration \( c = c_0 \) and a parabolic flow profile

\[ v = 6u_0 \frac{y}{H} \left[ 1 - \frac{y}{H} \right] e_z \]  

at the inlet, and vanishing axial gradients\n\[ \frac{\partial v}{\partial z} = 0 \]  

at the outlet. At the walls, we assume a no-slip condition \( v = 0 \), while the condition for \( c \) depends on the specific physical system. Lastly, we fix the pressure level to be zero at the lower corner of the outlet \( p(L, 0) = 0 \).

2.5 Dimensionless, effective transport coefficients

The dependency of the effective diffusion, mobility, viscosity and density on the particle concentration \( \phi \) is described by the correction factors \( \chi_D(c) \), \( \chi_\mu(c) \), \( \chi_\eta(c) \), and \( \chi_\rho(c) \), respectively, which can be interpreted as dimensionless effective transport coefficients. The suspension density is a weighted average of the water density \( \rho_0 \) and the particle density \( \rho_p \),

\[ \chi_\rho(\phi) = (1 - \phi) + \frac{\rho_p}{\rho_0} \phi. \]  

For bio-fluids, particles are often neutrally buoyant, so \( \rho_p \approx \rho_0 \) and \( \chi_\rho \approx 1 \). In a comprehensive numerical study of finite-sized, hard spheres suspended in an Newtonian fluid at low Reynolds numbers (the Stokes limit), the remaining transport coefficients \( \chi_D \), \( \chi_\mu \) and \( \chi_\eta \) were calculated by Ladd\(^2\)\(^1\) at 14 specific values of \( \phi \) between 0.001 and 0.45. For each of the coefficients \( \chi \), we fitted a fourth-order polynomial to the discrete set of values \( \gamma^{-1} \), while imposing the condition \( \chi(0)^{-1} = 1 \) to ensure agreement with the dilute-limit value,

\[ \chi_\rho(\phi)^{-1} = 1 + 2.476(2)\phi + 7.53(2)\phi^2 - 16.34(5)\phi^3 + 83.7(3)\phi^4 + \epsilon(\phi^5), \]  

(7a)

\[ \chi_D(\phi)^{-1} = 1 + 1.714(3)\phi + 6.906(13)\phi^2 - 17.05(3)\phi^3 + 61.2(1)\phi^4 + \epsilon(\phi^5), \]  

(7b)

\[ \chi_\mu(\phi)^{-1} = 1 + 5.53(3)\phi + 43.4(3)\phi^2 - 149.6(9)\phi^3 + 562(3)\phi^4 + \epsilon(\phi^5), \]  

(7c)

where the uncertainty of the last digit is given in the parentheses. To quantify the quality of the fits \( \chi_{\text{fit}} \) compared to the simulated data points \( \chi_{\text{data}} \), we calculated the relative difference \( \Delta \chi = \chi_{\text{data}} - \chi_{\text{fit}} \). We found the standard deviation of \( \Delta \chi \) to be 0.003, 0.0019 and 0.0058 for \( \chi_D \), \( \chi_\rho \), and \( \chi_\mu \), respectively. Thus, the fourth-order polynomials describe the data points well. The resulting fits from eqn (7) and the original data points\(^2\)\(^1\) are plotted in Fig. 2. Note that the mobility \( \chi_\mu \) is the most rapidly decreasing function of concentration \( \phi \).

We note that the dilute-limit relation \( \mu_0 = \frac{1}{6\pi\eta_0} \) and the definitions \( \mu = \chi_\mu \mu_0 \) and \( \eta^{-1} = \chi_\eta \eta_0^{-1} \) do not imply equality of \( \chi_\mu \) and \( \chi_\eta \). In fact, \( \chi_\mu < \chi_\eta \), which highlights the complex nature of the interaction problem at high densities.

2.6 Dimensionless parameters

The dynamics of the system can be characterized by three dimensionless numbers. One is the particle concentration \( \phi_0 \) which is a measure of the magnitude of the hydrodynamic interactions. The other two are related to the relative magnitude of the diffusion, advection and migration current densities in eqn (2), each characterized by the respective time scales \( \tau_{\text{diff}} = \frac{H^2}{8D_0} \), \( \tau_{\text{adv}} = \frac{H}{2\eta_0} \), and \( \tau_{\text{mig}} = \frac{H}{2\mu_0 F_{\text{mig}}} \). A small time scale
means that the corresponding process is fast and thus dominating. The importance of advection relative to diffusion and of migration relative to advection is characterized by the Péclet number and the transient Strouhal number, respectively,

\[ \text{Pé} = \frac{\tau_{\text{diff}}}{\tau_{\text{adv}}} \]

\[ \text{St} = \frac{\tau_{\text{adv}}}{\tau_{\text{mig}}} \]

\[ \phi_t = \frac{4}{3} \pi d^3 c_0 \]

3 Numerical implementation

Following Nielsen and Bruus,\textsuperscript{30} we implement the governing eqn (2)–(6) in the general weak-form PDE module of the finite element software COMSOL Multiphysics v. 5.0.\textsuperscript{31} To avoid spurious, numerically induced, negative values of \( c(r) \) near vanishing particle concentration, we replace \( c \) by the logarithmic variable \( s = \log(c/c_0) \). This variable can assume all real numbers, and it ensures that \( c = c_0 \exp(s) \) is always positive. We implement the weak form using Lagrangian test functions of first order for \( p \) and second order for \( s, v_x, \) and \( v_y \).

A particular numerical problem needs to be taken care of in dealing with the advection in our system. The typical average flow speed is \( u_0 = 200 \mu \text{m s}^{-1} \) in our channels of height \( H = 50 \mu \text{m} \), corresponding to a Péclet number \( \text{Pé} = 1.7 \times 10^4 \). To resolve this flow numerically, a prohibitively fine spatial mesh is needed. However, this problem can be circumvented by artificially increasing the particle diffusivity to \( \chi_P = 50 \) instead of using eqn (7b). This decreases the Péclet number to \( \text{Pé} = 2.3 \times 10^3 \), but since \( \text{Pé} \gg 1 \) for both the actual and the artificial diffusivity, advection still dominates in both cases, and as discussed in appendix A, the simulation results for the artificial diffusivity are accurate, but a coarser mesh suffices. In this highly advective regime \( \text{Pé} \gg 1 \), the detailed transport in the active area depends on the remaining parameters: the transient Strouhal number \( \text{St} \) and the inlet concentration \( c_0 \).

More details on the numerical implementation are given in appendix A.

4 Model results and discussion

To exemplify our continuum effective-medium model, we present two examples using the same channel geometry, but different particle migration mechanisms, namely magnetophoresis (MAP) and acoustophoresis (ACP). The magnetophoresis model is the one introduced by Mikkelsen and Bruus,\textsuperscript{23} but here extended to include the concentration dependence of the effective transport coefficients \( \chi \) introduced in section 2.5. This model primarily serves as a model validation. The acoustophoresis model is inspired by the many experimental works reported in Lab on a Chip.\textsuperscript{34,32–35}

4.1 Results for magnetophoresis (MAP)

As in Mikkelsen and Bruus,\textsuperscript{23} we consider the parallel-plate system sketched in Fig. 3(a) of height \( H = 50 \mu \text{m} \), length \( L = 350 \mu \text{m} \), and the boundary conditions of section 2.4. The particles are paramagnetic with magnetic susceptibility \( \chi_{\text{mag}} = 1 \) and diameter \( 2a = 2 \mu \text{m} \). A pair of long parallel thin wires is placed perpendicular to the \( xy \)-plane with their midpoint at \( r_0 = (x_0, y_0) = (250, 55) \mu \text{m} \), separated by \( d = dx \), and carrying electric currents \( \pm I \). The currents lead to a magnetic field, which induces a paramagnetic force \( F_{\text{map}} \) on each particle resulting in magnetophoresis. For a particle at position \( r \), the force is given by

\[ F_{\text{map}} = \frac{2}{\pi} \frac{\chi_{\text{mag}}}{\chi_{\text{em}}} d \left( \frac{r - r_0}{|r - r_0|^3} \right) \]

where \( \mu_{\text{em}}^0 \) is the vacuum permeability. Note how the rapid fifth-power decrease as a function of the wire distance \( |r - r_0| \), signals that \( F_{\text{map}} \) is a short-ranged local force. Dipole–dipole interactions are neglected. Using the expression for \( F_{\text{map}} \), we can determine the migration time \( \tau_{\text{mig}} \) for a particle.

![Fig. 3](image-url)
to migrate vertically from the channel center 
\((x, y) = \left(x_0, \frac{1}{2} H \right)\), directly below the wires, to the wall 
\((x, y) = \left(x_o, H \right)\). From this we infer the magnetic transient Strouhal
number \(St_{\text{map}}\),

\[
St_{\text{map}} = \frac{\Delta \phi_{\text{map}}}{\tau_{\text{map}}} = \frac{16 \mu J_{\text{map}}}{3 \pi^2 u_o \eta_o H^3}. \tag{10}
\]

Wall boundary conditions for the particles. Following the
simplifying model of Mikkelsen and Bruus, we assume that
once a particle reaches the wall near the wires, it sticks and
is thus removed from the suspension. For the upper wall at
\(y = H\), we therefore set the total particle flux in the y-direction
equal to the active migration \(J_y = cF_y\), while at the pas-
vive bottom wall at \(y = 0\), we choose the standard no-flux con-
dition \(J_y = 0\). For short times, this is a good model, but event-
tually in a real system, the build-up of particles at the wall is
likely to occur, which leads to a dynamically change of the
channel geometry.

Particle capture as a function of \(\phi_0\) and \(St_{\text{map}}\). In Fig. 4 we
have plotted the normalized steady-state particle concentra-
tion \(\phi(r)/\phi_0\) of the system shown in Fig. 3(a) for the nine combi-
nations of \(\phi_0 = 0.001, 0.01, 0.1\) with \(St_{\text{map}} = 1, 2, 5\). Firstly,
it is evident from the figure that for a given fixed \(\phi_0\) and in-
creasing \(St_{\text{map}}\), an increasing number of particles are cap-
tured by the wire-pair, thus decreasing \(\phi(r)\) in a region of in-
creasing size down-stream from the wires. Note also how the
slow-moving particles near the bottom wall are pulled up-
wards by the action of the wires, leaving a decreased particle
density there. Secondly, for a fixed \(St_{\text{map}}\), the particle capture
is non-monotonic as a function of \(\phi_0\): it increases going from
\(\phi_0 = 0\) via 0.001 to 0.01, but then decreases going from 0.01
to 0.1. Thirdly, for the highest concentration \(\phi_0 = 0.1\), the
morphology of the depleted region has changed markedly: a
small region depleted of particles forms around the wires at the
top wall together with narrow downstream depletion re-
gions along both the top and bottom wall. All three regions
increase somewhat in size for increasing \(St_{\text{map}}\).

The normalized capture rate \(\beta_{\text{map}}\). We analyze the particle
capture quantitatively in terms of the normalized capture rate
\(\beta_{\text{map}}\) defined as the ratio between the flux \(\Gamma_{\text{map}}\) of particles
captured on the top wall by the wires, and the flux \(\Gamma_{\text{inlet}}\) of
particles entering through the inlet,

\[
\beta_{\text{map}} = \frac{\Gamma_{\text{map}}}{\Gamma_{\text{inlet}}}, \tag{11a}
\]

\[
\Gamma_{\text{surf}} = \int_n d\mathbf{J}. \tag{11b}
\]

where \(n\) is the surface normal of a given surface. In Fig. 5(a),
we have plotted \(\beta_{\text{map}}\) as a function of \(\phi_0\) for different \(St_{\text{map}}\)
values. All of the curves in this plot show the same qualitative
behavior: the capture fraction increases with inlet concentra-
tion until a certain point, beyond which it declines.

This general behavior is illustrated by the three specific
cases \(\phi_0 = 0.02, 0.045, \) and 0.1 denoted by the purple bullets
A, B and C on the purple “\(St_{\text{map}} = 1\)”-curve. For these three
cases, we have in Fig. 5(b1–d1) plotted the magnetically-
induced change \(\Delta v = v - v_0\) in the suspension velocity, and in
Fig. 5(b2–d2) the magnetic force density \(f = cF_{\text{map}}\) overlaid by con-
tour lines of \(c\) (white). For A in Fig. 5(b1) and (b2) with
\(\phi_0 = 0.02\), more particles are captured compared to the dilute
case, because \(cF_{\text{map}}\) induces a clockwise flow-roll that advects
the upstream particles closer to the wires. For B in
Fig. 5(c1) and (c2) with \(\phi_0 = 0.045\), \(cF_{\text{map}}\) induces two flow
rolls, of which the upstream one rotates counter clockwise
and advects particles away from the wires, thus lowering the
capture compared to A. For C in Fig. 5(d1) and (d2) with \(\phi_0 = 0.1\), \(cF_{\text{map}}\) induces a single strong counter-clockwise flow roll
that more efficiently advects the upstream particles away
from the wires, and \(\beta_{\text{map}}\) decreases further.

While the initial increase in \(\beta_{\text{map}}\) for \(\phi_0\) increasing from
0 to 0.03 is in agreement with Mikkelsen and Bruus, black
circles in Fig. 5(a), the following decrease in \(\beta_{\text{map}}\) for \(\phi_0\)

![Fig. 4](image-url)

Fig. 4 Heat color plot of the normalized concentration field \(\phi(r)/\phi_0\) from 1 (white, inlet concentration) to 0 (black, no particles) in the
magnetophoretic system Fig. 3(a), for the nine combinations of \(\phi_0 = 0.001, 0.01, 0.1\) with \(St_{\text{map}} = 1, 2, 5\).
increasing beyond 0.03 was not observed by them. The deviation is not due to our inclusion of the effective transport coefficients \( \chi_\eta \) and \( \chi_\mu \), since repeating the simulations\(^{23} \) with \( \chi_\eta = \chi_\mu = 1 \) resulted in a curve indistinguishable from the green curve \( St_{\text{map}} = 0.7 \) in Fig. 5(a). We speculate that this deviation is due to our improved mesh resolution.

### 4.2 Results for acoustophoresis (ACP)

We consider a system identical to the one described in section 4.1, except that the single-particle migration force \( F_{\text{mig}} \) now is taken to be the acoustic radiation force \( F_{\text{acp}} \) arising from the scattering of an imposed standing ultrasound wave on the microparticles. It is given by\(^{36} \)

\[
F_{\text{acp}} = - \nabla U_{\text{ac}}, \tag{12a}
\]

where \( U_{\text{ac}} = \frac{4\pi}{3} a^3 \left[ \text{Re}[f_0] - \frac{1}{2} \text{Re}[f_1] \langle \nu_0^2 \rangle - \frac{3}{4} \text{Re}[f_1] \langle \nu_1^2 \rangle \right] \).\(^{12b} \)

The acoustic potential \( U_{\text{ac}} \) is obtained by inserting the acoustic fields (13) into eqn (12), and the result is a force, which focuses the particles at the pressure node in the channel center at \( y = \frac{1}{2} H \), see appendix B for details.

Given \( F_{\text{acp}} \), we then calculate the transient Strouhal number from the acoustic migration time \( \tau_{\text{mig}} \) in the dilute limit \( \chi_\mu = 1 \) for a single particle migrating from \( (x, y) = \left( \frac{1}{2}, \frac{1}{8} H \right) \) to \( (x, y) = \left( \frac{1}{2}, \frac{3}{8} H \right) \)\(^{26} \) and from the advection time given by \( \tau_{\text{adv}} = \frac{H}{4u_0} \),

\[
St_{\text{acp}} = \frac{\tau_{\text{adv}}}{\tau_{\text{mig}}} = \frac{4\pi^2}{\ln \left( \frac{\text{col} \pi}{8} \right)} \frac{E_{\text{ac}} \Phi}{\eta_0 H u_0}. \tag{14}
\]

Here, \( E_{\text{ac}} = p_a^2 (4\rho_0 c_{\text{liq}}^2) \) is the acoustic energy density (typically 1–100 \( \text{J m}^{-3} \)), \( c_{\text{liq}} \) is the speed of sound in the carrier liquid, and \( \Phi = \frac{1}{3} f_0 + \frac{1}{2} f_1 \) is the so-called acoustic contrast factor.\(^{37} \)

Wall boundary condition for the particles. As the particles in this model enter from the inlet, focus towards the channel center \( y = \frac{1}{2} H \), and leave through the outlet, we choose the simple no-flux boundary condition, \( n f_1 |_{y=0,H} = 0 \), for the bottom and top walls.

Particle capture as a function of \( \phi_0 \) and \( St_{\text{acp}} \). In Fig. 6 is plotted the normalized steady-state particle concentration

\[
\begin{align}
\beta_{\text{map}} &= \frac{\text{Re}[f_0]}{\text{Re}[f_0] - \frac{1}{2} \text{Re}[f_1] \langle \nu_0^2 \rangle - \frac{3}{4} \text{Re}[f_1] \langle \nu_1^2 \rangle}. \tag{12b}
\end{align}
\]

Here, \( \omega = 2\pi f \) is the angular actuation frequency with the frequency \( f \) typically in the MHz range, and \( p_a \) is the acoustic pressure amplitude typically in the 0.1 MPa range.\(^{26,32} \) The acoustophoretic force \( F_{\text{acp}} \) is obtained by inserting the acoustic fields (13) into eqn (12), and the result is a force, which focuses the particles at the pressure node in the channel center at \( y = \frac{1}{2} H \), see appendix B for details.
of the system shown in Fig. 3(b) for the nine combinations of $\phi_0 = 0.001, 0.01, 0.1$ to $St_{acp} = 0.5, 1, 2$. Clearly, for a given fixed $\phi_0$ and increasing $St_{acp}$, an increasing number of particles are focused in the center region. This is expected since $St_{acp}$ is proportional to the acoustic energy density $E_{ac}$. For $\phi_0 = 0.001$, the center-region concentration (white) is 1.73 times the inlet concentration (orange). Secondly, for a fixed $St_{acp}$, the particle focusing is monotonically decreasing as a function of $\phi_0$ as seen from the decrease in both the width of the particle-free black regions and the magnitude of the downstream center-region concentration. See in particularly the change from $\phi_0 = 0.01$ to 0.1.

To quantify the acoustic focusing for a given inlet concentration, we introduce in analogy with eqn (11) the normalized flux $\beta_{acp}$ of focused particles leaving through the middle 30% of the outlet,

$$\beta_{acp} = \frac{G_{center}}{G_{inlet}}$$ (15a)

$$G_{center} = \int_{0.55HF}^{0.45HF} dy J_y(L,y),$$ (15b)

In the lin–log plot Fig. 7(a) of $\beta_{acp}$ versus $\phi_0$ for the four Strouhal numbers $St_{acp} = 0.1, 0.5, 1, 2$, the above-mentioned monotonic decrease in acoustophoretic focusing is clear. Especially for $\phi_0$ exceeding 0.01 a dramatic reduction in focusability sets in.

To determine the dominant hydrodynamic interaction mechanism behind this reduction, we introduce the normalized particle focusing fraction $\Delta\beta_{acp}(\phi_0, St_{acp})$ by the definition,

$$\Delta\beta_{acp}(\phi_0, St_{acp}) = \frac{\beta_{acp}(\phi_0, St_{acp}) - \beta_{acp}(\phi_0, 0)}{\beta_{acp}(0, St_{acp}) - \beta_{acp}(0, 0)}.$$ (16)

Here, the difference in normalized center out-flux $\beta_{acp}$ at an inlet concentration $\phi_0$ between “acoustics on” ($St_{acp} > 0$) and “acoustics off” ($St_{acp} = 0$), is calculated relative to the same difference computed in the dilute limit $c_{mig} = 0$ and $\chi_\mu = \chi_\eta = 1$, which is written as $\phi_0 = 0$. We notice that $\Delta\beta_{acp}(\phi_0, St_{acp})$ is unity if there is no hydrodynamics interaction effects, and zero if there is no acoustophoretic focusing. In Fig. 7(b), using the full model ($\chi_\mu = \chi_\mu(\phi)$ and $\chi_\eta = \chi_\eta(\phi)$ denoted “all effects”), $\Delta\beta_{acp}$ is plotted versus $\phi_0$ (black lines) for the same four $St_{acp}$ values.
as in Fig. 7(a). We note that for the full model (black lines), $\Delta \phi_{\text{map}}$ is almost independent of the strength $St_{\text{ap}}$ of the acoustophoretic force, which makes it ideal for studying the interaction effects through its $\phi_0$ dependency.

To elucidate the origin of the observed reduction in focusability, we also plot the following special cases: “no effects” with $\chi_\mu = \chi_\eta = 1$ and $\phi_{\text{mig}} = 0$ (red lines), “only $\phi_{\text{mig}}$” with $\chi_\mu = \chi_\eta = 1$ (blue lines), “only $\chi_\eta$” with $\chi_\mu = 1$ and $\chi_\eta = \chi_\eta(\phi)$ (purple lines), and “only $\chi_\mu$” with $\chi_\mu = \chi_\mu(\phi)$ and $\chi_\eta = 1$ (green lines). These curves reveal that most of the interaction effect is due to the effective mobility: the green curves “only $\chi_\eta$” are nearly identical to the black curves “all effects”, and both groups of curves show no or little dependence on $St_{\text{ap}}$.

The effective suspension viscosity $\chi_\eta(\phi)$ and the particle-induced bulk force $\phi_{\text{mig}}$ on the suspension only have a minor influence on $\Delta \phi_{\text{map}}$.

Finally, in Fig. 8 we show the first step towards an experimental validation of our continuum model of hydrodynamic interaction effects. In the figure are seen two micrographs of acoustophoretic focusing of 5.1 $\mu$m-diameter polystyrene particles in water in the silicon-glass microchannel setup described in Augustsson et al.\textsuperscript{26} The micrographs are recorded at Lund University by Kevin Cushing and Thomas Laurell under stop-flow conditions $u_0 = 0$ after 0.43 s of acoustophoresis with the acoustic energy $E_{\text{ac}} = 10^2$ J m$^{-3}$. The field of view is near the right edge $x_2$ of the active region sketched in Fig. 3(b).

The good focusing obtained at the low density $\phi_0 = 0.001$ is well reproduced by our model, see Fig. 8(a) and (c). Likewise, the much reduced focusing obtained at the high density $\phi_0 = 0.1$ is also well reproduced by our model, see Fig. 8(b) and (d). Moreover, which is not shown in the figure, the observed transient behavior of the two cases is also reproduced by simulations: For the low concentration $\phi_0 = 0.001$, the particles are migrating towards to pressure node directly in straight lines, while for the high concentration $\phi_0 = 0.1$, pronounced flow rolls are induced, similar to those shown in Fig. 5.

### 4.3 Discussion

The results of the continuum effective model simulations of the MAP and the ACP models show that the dominating transport mechanisms differ. For the ACP model, the behavior of a suspension in the high-density regime $\phi_0 \gtrsim 0.01$ is dominated by the strong decrease in the effective single-particle mobility $\chi_\mu$ shown in Fig. 2. Here, the particle–particle interactions are so strong that the suspension moves as a whole with suppressed relative single-particle motion. The response is a steady degradation of the focusing fraction $\beta_{\text{ap}}$ shown in Fig. 7, as it becomes increasingly difficult for a given particle to migrate towards the center line, where the concentration increases and the mobility consequently decreases. In contrast, for the MAP model, the capture fraction $\beta_{\text{map}}$ increases as $\phi_0$ increases from 0 to 0.03, confirming the results of Mikkelsen and Bruus.\textsuperscript{23} The reason for the increased capture is that the particle concentration in the near-wire region is low due to the removal of particles, leaving the single-particle mobility unreduced there, while in the region a little farther away, particles attracted by the wires pull along other particles in a positive feedback due to the high-density “stiffening” of the suspension there. Thus, for the MAP system, the dominating interaction effect is through the particle-induced body force $\phi_{\text{mig}}$ in eqn (5a).

For $\phi_0$ increasing beyond 0.03, the observed non-monotonic behavior of $\beta_{\text{map}}$ and continued monotonic behavior of $\beta_{\text{ap}}$ as a function of increasing $\phi_0$, depends on whether or not the specific forces and model geometry lead to the formation of flow rolls, such as shown for the MAP model in Fig. 5. The flow-roll-induced decrease observed in the capture efficiency $\beta_{\text{map}}$ for $\phi_0 \gtrsim 0.03$, see Fig. 7(a), was not observed by Mikkelsen and Bruus,\textsuperscript{23} presumably because of their relatively coarse mesh. For the ACP model, no strong influence from induced flow rolls was observed, and the resulting decrease in focus efficiency was monotonic as a function of $\phi_0$. This difference in response is caused by differences in spatial symmetry and in the structure of the two migration forces. The MAP model is asymmetric around the horizontal center line, while the ACP model is symmetric, which accentuates the appearance of flow rolls in the former case. Moreover, the rotation of the body force density $\phi_{\text{mig}}$ can be calculated as $\nabla \times [\phi_{\text{mig}}] = \nabla \times \phi_{\text{mig}} - \phi_{\text{mig}} \times \nabla c$. In the MAP model, both of these rotations are non-zero, while in the ACP model, the first term is zero as $F_{\text{ap}}$ is a gradient force, and this feature leads to significant flow rolls in the MAP model and weak ones in the ACP model.

In the above discussion of the results of our model, we have attempted to describe general aspects of hydrodynamic interaction effects, which may be used in future device design considerations. Given the sensitivity on the specific...
5 Conclusion

We have presented a continuum effective medium model including hydrodynamic particle–particle interactions, which requires only modest computational resources and time. The results presented in this paper could be obtained in the order of minutes on an ordinary PC. We have validated the model by comparison to previously published models and found quantitative agreement. Moreover, as a first step towards experimental validation, we have found qualitative agreement with experimental observations of acoustophoretic focusing.

Using our numerical model, we have shown that the response of a suspension to a given external migration force depends on the system geometry, the flow, and the specific form of the migration forces. At high densities \( \phi_0 \gtrsim 0.01 \) the suspension “stiffens” and relative single-particle motion is suppressed due to a decreasing effective mobility \( \chi_\mu \). However, the detailed response depends on whether or not particle-induced flow rolls appear with sufficient strength to have an impact.

The continuum effective medium model may prove to be a useful and computationally cheap alternative to the vastly more cumbersome direct numerical simulations in the design of lab-on-a-chip systems involving high-density suspensions, for which hydrodynamic particle–particle interactions are expected to play a dominant role.

Appendix A: Numerical implementation, details

The numerical implementation has been carefully checked and validated. To avoid spurious numerical errors, we have performed a mesh convergence analysis as described in details in Muller et al.\(^{37} \) For a continuous field \( g \) on a given mesh, we have calculated the relative error

\[
C(g) = \sqrt{\int [g - g_{rel}]^2 \, da / \int [g_{rel}]^2 \, da},
\]

where \( g_{rel} \) is the solution obtained for the finest possible mesh. For all fields involved in this study, \( C \) falls in the range between \( 10^{-3} \) and \( 10^{-3} \), and it has the desired exponential decrease as function of \( h_{\text{mesh}}^{-1} \), where \( h_{\text{mesh}} \) is the mesh element size. The number of mesh points and degrees of freedom used in our final simulations are listed in Table 1.

Table 1: The number of mesh elements \( N_{\text{mesh}} \), the degrees of freedom DOF, and the mesh Pécelt number \( P\epsilon_{\text{mesh}} \) used in the numerical simulations of the MAP and ACP models

<table>
<thead>
<tr>
<th>Model</th>
<th>( N_{\text{mesh}} )</th>
<th>DOFs</th>
<th>( P\epsilon_{\text{mesh}} )</th>
</tr>
</thead>
<tbody>
<tr>
<td>MAP</td>
<td>70 000–225 000</td>
<td>200 000–2 000 000</td>
<td>5–10</td>
</tr>
<tr>
<td>ACP</td>
<td>50 000–200 000</td>
<td>400 000–2 000 000</td>
<td>5–10</td>
</tr>
</tbody>
</table>

We have compared our results with previous work in the literature. Setting \( \chi_D = \chi_\eta = 1 \), our MAP model in section 4.1 is identical to the one by Mikkelsen and Bruus.\(^{23} \) For this setting, we have calculated the capture fraction \( \beta_{\text{map}} \) as a function of the rescaled quantity \( (D\bar{\gamma})^2/\bar{u}_0 \) (which also appears in the definition (10) for \( St_{\text{map}} \)) and found exact agreement, including the remarkable data collapse, with the results presented in Fig. 3 of Mikkelsen and Bruus.\(^{23} \)

The numerical stability of the finite-element simulation is controlled by the mesh Pécelt number \( P\epsilon_{\text{mesh}} = h_{\text{mesh}} \mu_0 \sqrt{\chi_\mu} \), that involves the mesh size \( h_{\text{mesh}} \) instead of a physical length scale. We found that convergence was ensured for \( P\epsilon_{\text{mesh}} \lesssim 10 \), which implies the condition \( h_{\text{mesh}} \lesssim 10\bar{\gamma}_0 \bar{D}/\bar{u}_0 \) for the mesh size. The higher an advection velocity \( \bar{u}_0 \), the finer a mesh is needed. In order to obtain computation times of the order of minutes per variables (\( \bar{P}e, S\bar{t}, \phi_0 \)) on our standard PC, while still keeping \( P\epsilon_{\text{mesh}} \lesssim 10 \), we artificially increased \( \chi_D \). In Fig. 9 we have plotted the capture fraction \( \beta_{\text{map}} \) and the focusing fraction \( \beta_{\text{ACP}} \) as functions of the transient Strouhal number \( S\bar{t} \) for \( \chi_D \) ranging from 25 to 10 000 in the dilute limit with \( \chi_\mu = \chi_\eta = 1 \) and \( cF_{\text{mig}} = 0 \). It is clear from Fig. 9 that the \( \beta(S\bar{t}) \)-curves converge for sufficiently small values of \( \chi_D \). For \( \chi_D \lesssim 50 \), we determine that acceptable convergence is obtained in both models, and we thus fix \( \chi_D = 50 \) for all calculation performed in this work. This choice of \( \chi_D = 50 \) allows for acceptable computation times (~10 min per variable set). In other words, once the Pécelt number is sufficiently high, the convergence is ensured by the dominance of advection over diffusion.
Appendix B: The acoustophoretic force

The acoustophoretic force $F_{ac}$ is found from the gradient of $U_{ac}$ in eqn (12) using the acoustic fields $p_1$ and $v_1$ defined in eqn (13). In Fig. 10 is shown a combined color and vector plot of the resulting force $F_{ac}$ given by the explicit expressions of its components $F_{ac,x}$ and $F_{ac,y}$:

\[
F_{ac,x} = \frac{4}{3} \pi \alpha^2 E_{ac} \kappa_2 \sin(2\tilde{x}) \times \left[ \left( f_0 + \frac{3k^2}{2k^2} f_1 \right) \cos^2 \tilde{y} + \frac{3k^2}{2k^2} f_1 \sin^2 \tilde{y} \right]},
\]

(17a)

\[
F_{ac,y} = -\frac{4}{3} \pi \alpha^2 E_{ac} \kappa_2 \sin(2\tilde{y}) \times \left[ \left( f_0 + \frac{3k^2}{2k^2} f_1 \right) \sin^2 \tilde{x} - \frac{3k^2}{2k^2} f_1 \cos^2 \tilde{y} \right].
\]

(17b)

Here, $\tilde{x} = k_x \left( x - \frac{2}{7} \right)$ and $\tilde{y} = k_y y$ are dimensionless coordinates, and $k_x = \frac{\pi}{4H}$, $k_y = \frac{\pi}{H}$, and $k^2 = k_x^2 + k_y^2$ are wave numbers. For water-suspended polystyrene spheres with radius $\alpha = 1 \mu$m, we have $f_0 = 0.444$ and $f_1 = 0.034$. The acoustic energy density is set to $E_{ac} = 10^2$ J m$^{-3}$.

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References


