In a recent experiment, Segrè et al. [1] used the particle imaging velocimetry (PIV) technique to measure the spatial correlation function, \( C(\ell) = \langle \delta v(r) \delta v(r + \ell) \rangle \), of the velocity fluctuation \( \delta v \) in a sedimenting suspension of non-Brownian particles over a wide range of particle concentrations and sample sizes. They found that the measured \( C(\ell) \sim \exp(-\ell/\xi) \), where the velocity correlation length \( \xi \) depends on the particle radius \( a \) and volume fraction \( \phi_0 \) in a non-trivial power-law form \( \xi \approx a \phi_0^{-1/3} \). In this paper we propose a new set of coarse-grained equations of motion to describe concentration and velocity fluctuations in a dilute sedimenting suspension of non-Brownian particles. With these equations we find that colloidal sedimentation is analogous to high Rayleigh number, high Prandtl number turbulent convection [2,3]. Our model explains the experimental results by Segrè et al. and also provides a coherent framework for the study of sedimentation dynamics in different colloidal systems.

To understand the basic principles governing the colloidal sedimentation, we consider a simple case of a dilute sedimenting suspension of hard spheres in a long cylindrical tube of radius \( L \). To separate the velocity fluctuation \( \delta u \) from the mean settling velocity \( \bar{u} \), we choose a uniform suspension with \( \bar{u} = 0 \) as our reference system. It has been suggested [1,4] that velocity fluctuations in a sedimenting suspension may arise from fluctuations of the local particle concentration. Therefore, we model the colloidal sedimentation with a coarse-grained Navier-Stokes equation. The fluid velocity \( \delta u \) and pressure \( \delta p \) at a point \( x \) satisfy the creeping flow equation [5]

\[
\nabla \delta p(x) - \eta \nabla^2 \delta u(x) = f \delta u(x),
\]

where \( \eta \) is the viscosity of the fluid and \( f \delta u \) represents the fluctuation of the particle number density \( n(x) \) about its mean \( \bar{n} \). In the above, \( f = (4 \pi/3) \alpha^2 \Delta \rho g \) is the buoyancy force acting on a particle of radius \( a \), where \( g \) is the gravitational acceleration and \( \Delta \rho = \rho_p - \rho_s \) is the density difference between the particle \( \rho_p \) and the solvent \( \rho_s \). In writing Eq. (1) we have assumed that the fluid volume element \( \delta V \) is a coarse-grained volume, which is large enough to contain many particles but is small enough such that the particle distribution inside \( \delta V \) is uniform. In this case, we have \( f \delta n(x) = \Delta \rho g [\phi(x) - \phi_0] \), where \( \phi(x) \) is the particle volume fraction and \( \phi_0 \) is its mean value.

Nondimensionalizing Eq. (1) with respect to the length \( L \), the time \( L^2/\eta \), and the concentration \( \phi_0 \), we have

\[
- \frac{1}{\sigma} \nabla \delta p(x) + \nabla^2 \delta u(x) = Ra \phi(x) \delta z,
\]

where the unit vector \( \delta z \) is directed upward opposite to the direction of \( g \), and the dynamic pressure \( \delta p \) has included a term, \(-\Delta \rho g \phi_0 \), to absorb contributions from the constant forcing term \(-\Delta \rho g \phi_0 \). In Eq. (2) the Rayleigh number \( Ra \) is defined as

\[
Ra = \Delta \rho g \phi_0 L^2 / (\eta \Delta p),
\]

where \( D \) is an effective diffusion constant of the particles. The Schmidt number \( \sigma \) is given by \( \sigma = \nu / \eta \) with \( \nu \) being the kinematic viscosity of the fluid. For a dilute suspension of small colloidal particles, \( D \) is approximately equal to the particle self diffusion constant \( D_p = k_B T/(6 \pi \eta a) \), where \( k_B T \) is the thermal energy. For large non-Brownian particles, however, the effect of thermal agitations is negligible and their diffusion-like motion is produced by the hydrodynamic interactions between the particles [6]. Nicolai et al. have shown [7] that the hydrodynamic diffusivity has the form \( D_h \sim 5 a^2 \eta \), where \( a^2 \eta \) is the Stokes velocity.

With the hydrodynamic diffusivity \( D_h \), Eq. (3) becomes

\[
Ra = 0.9 \phi_0 (L/a)^3.
\]

It should be mentioned that while it is cancelled out in \( Ra \), \( \Delta \rho g \) is needed so that \( D_h \) can be used to describe the hydrodynamic diffusion of the settling particles at small length scales. Equation (2) together with the continuity equation for an incompressible fluid

\[
\nabla \cdot \delta u = 0
\]

and the advective mass diffusion equation

\[
\partial_t \phi + (\delta u \cdot \nabla) \phi = \nabla^2 \phi
\]

complete the description of concentration and velocity fluctuations in colloidal sedimentation.

It is evident that Eqs. (2)-(6) are the same as those for buoyancy-driven convection [3]. Velocity and concentration fluctuations in colloidal sedimentation are therefore analogous to those in buoyancy-driven convection, and they are completely controlled by the two dimensionless parameters \( Ra \) and \( \sigma \), once the boundary conditions are specified. We now estimate typical values of \( Ra \) and \( \sigma \) in colloidal sedimentation.

In the experiment by Segrè et al. [1], the particle’s radius \( a \approx 8 \mu m \), Stokes velocity \( U_0 \approx 6.5 \mu m/s \), volume fraction \( \phi_0 \approx 0.05 \), and the characteristic sample size \( L \approx 1 \text{ cm} \). With these experimental values, we find \( D_h \approx 2.6 \times 10^{-6} \text{ cm}^2/\text{s} \), \( Ra \approx 8.8 \times 10^7 \), and \( \sigma \approx 3800 \). The Schmidt number \( \sigma \)
is equivalent to the Prandtl number in thermal convection. Colloidal sedimentation is, therefore, associated with high Rayleigh number, high Prandtl number turbulent convection.

To understand the sedimentation dynamics, it is helpful to distinguish two characteristic length scales in convection: the viscous dissipation length \( \delta_v \), and the diffusive dissipation length \( \delta_d \). The values of \( \delta_v \) and \( \delta_d \) are determined, respectively, by the transition Reynolds number \( Re_t = \delta_u \delta_d / \nu \) and the transition Peclet number \( Pe_c = \delta_u \delta_d / D_h \). Here \( \delta_u \) is the rms value of the velocity fluctuation \( \delta u \) averaged over a volume of \( \delta_v^3 \) (or \( \delta_d^3 \)). It is the ratios of these lengths to each other and to the sample size \( L \) that determine the flow state of the system [2]. For high-\( R_a \), high-\( \sigma \) turbulent convection, one anticipates that the flow consists of three different regions: (i) \( a < \ell < \delta_v \), (ii) \( \delta_v < \ell < \delta_d \), and (iii) \( \delta_d < \ell < L \). In Region (i), molecular viscosity and hydrodynamic diffusivity determine the momentum and mass transport processes, respectively, and hence the particle distribution remains uniform without any large fluctuations. In Region (ii) turbulent (or eddy) diffusivity and molecular viscosity are dominant, and thus large fluctuations in particle concentration are expected but the velocity field remains relatively smooth. Finally, in Region (iii) turbulent diffusivity and viscosity both dominate over the corresponding hydrodynamic and molecular processes. In this case, one expects to see large fluctuations both in particle concentration and in velocity at different length scales.

We first discuss the length \( \delta_d \) above which velocities become large and concentration fluctuations are transported by convection. This occurs when the local Peclet number \( Pe_c = \delta_u \ell / D_h \) becomes larger than \( Pe_c \). Recent thermal convection experiments have shown [8] that while turbulent mixing creates on average an isothermal fluid in the turbulent bulk region, large temperature fluctuations still remain in the region and the characteristic length scale associated with these fluctuations is of the order of \( \delta_v \). Therefore, the velocity correlation length \( \xi \) is determined by \( \delta_d \) in Region (ii). According to Kraichnan's theory [2],

\[
\delta_d \simeq \left( \frac{2 \pi^2 Pe_c^2}{L R_a} \right)^{1/3},
\]

where the power law amplitude is expressed in terms of the numerical value of \( Pe_c \). Priestley [9] first gave a direct argument for the \( R_a^{-1/3} \) scaling. He argued that when \( R_a \) is large enough, \( \delta_d \) should be a new length scale independent of the sample size \( L \). With Eqs. (7) and (4), we immediately have \( \xi \sim \delta_d \sim LR_a^{-1/3} \sim a \delta_0^{-1/3} \). The mapping of colloidal sedimentation to turbulent convection, therefore, explains the experimental finding that \( \xi \sim 11 a \delta_0^{-1/3} \). It also provides a physical interpretation for the existence of a velocity cut-off length, which prevents hydrodynamic dispersion coefficients from being divergent.

We now discuss the velocity variance \( \delta u \) in colloidal sedimentation. According to Kraichnan's theory [2],

\[
\delta u \simeq \frac{Pe_c D_h}{\delta_d} \frac{Pe_c D_h}{\left( 2 \pi^2 Pe_c^2 \right)^{1/3} L R_a^{-1/3}}.
\]

Eq. (8) states that at the transition Peclet number \( Pe_c \), the mass flux due to hydrodynamic diffusion, \( D_h \delta u / \delta_d \), is approximately equal to that by convection, \( \delta u \delta_p \). Because \( \delta_d \simeq \delta_0^{1/3} \) and \( D_h \simeq a \delta_0 \), we find from Eq. (8) that \( \delta u \sim \delta u \simeq D_h / \delta_d \simeq \delta_0^{1/3} \), which is independent of the sample size \( L \). This result agrees well with the experimental finding that \( \delta u \sim 2 \delta_0^{1/3} \) [1].

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