

Hydrodynamics

Contents

1 Objectives	3
1.1 Objetivos específicos conceptuales:	3
1.2 Objetivos específicos procedimentales:	3
1.3 Objetivos específicos actitudinales:	3
2 Continuum dynamics	4
2.1 The velocity field	4
2.1.1 Streamlines	4
2.1.2 Particle Trajectories	5
2.1.3 Streaklines	5
2.2 Mass conservation	5
2.3 Incompressible flow	6
2.3.1 Global and local forms	6
2.3.2 Leonardo's Law	7
2.4 Lagrangian and Eulerian derivatives	7
2.5 Cauchy's equation	8
3 Incompressible frictionless (Euler) fluids	9
3.1 Euler equation for incompressible ideal flow	9
3.2 Adiabatic and isentropic flow	10
3.2.1 Noundary conditions	11
3.2.2 irrotational flow	11
3.3 Bernoulli's equation	12
3.3.1 The Venturi effect	13
3.3.2 Torricelli's law	13
3.3.3 Quasi-stationary emptying	14
3.3.4 Case: The Pitot tube	15
3.3.5 Rotating fluid in steady state	15
3.4 Vorticity	16
3.4.1 The vorticity field	16
3.4.2 Vortex lines	16
3.4.3 Equation of motion for vorticity	17
3.5 Circulation	17
3.5.1 Stokes' theorem	17
3.6 Potential flow	18
3.6.1 Incompressible potential flow	18
3.7 Application: Cylinder in uniform crosswind	19
3.7.1 Velocity potential	19
3.7.2 Velocity and pressure fields	19
3.7.3 Lift on a half-cylinder	20

4 Viscosity	20
4.1 Shear viscosity	20
4.1.1 Molecular origin of viscosity in gases	21
4.1.2 Temperature dependence of viscosity	21
4.1.3 Kinematic viscosity	21
4.1.4 Velocity-driven planar flow	22
4.1.5 Case: Steady planar flow	22
4.1.6 Case: Viscous friction	22
4.1.7 Case: Shear wave (Stokes' second problem)	23
4.1.8 Case: Free momentum diffusion	23
4.1.9 Case: Growth of a boundary layer (Stokes' first problem)	23
4.2 Dynamics of incompressible Newtonian fluids	24
4.2.1 Isotropic viscous stress	24
4.3 The Navier–Stokes equations	25
4.4 Boundary conditions	25
4.5 Classification of flows	26
4.5.1 The Reynolds number	26
4.5.2 Examples	26
4.5.3 Hydrodynamic similarity	27
4.5.4 Example: Robofly	28
4.5.5 Example: High-pressure wind tunnels	28

Disclaimer

Discussions taken from Alonso Sepúlveda [1], Landau & Lifschits [2], Lautrup [3] books

1 Objectives

1.1 Objetivos específicos conceptuales:

- OC12. Enumerar los campos que se utilizan típicamente para describir fluidos en movimiento (campos básicos de la hidrodinámica).
- OC13. Reconocer las condiciones necesarias para aproximar el movimiento de un fluido como: incompresible, irrotacional, no viscoso.
- OC14. Distinguir la descripción euleriana y lagrangiana de un fluído.
- OC15. Definir la viscosidad tanto desde el punto de vista microscópico como del punto de vista macroscópico.
- OC16. Definir los conceptos de: vorticidad, potencial de velocidades, campo de corriente.

1.2 Objetivos específicos procedimentales:

- OP10. Deducir las ecuaciones de movimiento de un fluido, mediante la aplicación de la descripción Lagrangiana y la segunda ley de Newton.
- OP11. Enunciar y demostrar el teorema de Bernoulli y estudiar sus aplicaciones.
- OP12. Deducir el tensor de viscosidad para un fluido Newtoniano isotrópico e identificar de allí las viscosidades de corte y volumétrica.
- OP13. Deducir las ecuaciones de Navier-Stokes para fluidos Newtonianos.
- OP14. Resolver analíticamente algunas situaciones hidrodinámicas simples.
- OP15. Discretizar las ecuaciones de Navier-Stokes en situaciones simples (flujo impulsado por presión, flujo impulsado por velocidad, flujo estacionario alrededor de un obstáculo) y resolverlas numéricamente.

1.3 Objetivos específicos actitudinales:

- OA4. Reconocer la importancia de los métodos numéricos en la solución de problemas de hidrodinámica.

2 Continuum dynamics

Continuous matter in motion appears throughout nature: flowing water, atmospheric currents, and even the vibrations of solids such as church bells, trees, or the ground during earthquakes. Although these phenomena are diverse, they are all governed by Newton's equations of motion applied to continuous media. While these equations are easy to formulate, analytic solutions typically exist only in highly idealized situations, and deeper understanding often requires experiments or numerical simulation.

Solids differ from fluids in that they resist changes of shape and tend to retain their form, whereas fluids flow and do not. Both, however, are treated within the same framework of continuum mechanics. Two fundamental principles govern their motion: conservation of mass and Newton's Second Law (momentum balance) applied to a continuous system. With appropriate expressions for internal and external forces, these lead to the equations of motion for the mass density and velocity fields.

In this chapter the analysis will focus on continuous matter in motion, with emphasis on fluids. The same equations of motion will also be applied on cosmological scales, showing that they yield a surprisingly coherent description of the universe. Later chapters will address more terrestrial applications of matter in motion.

The mathematical description of the state of a moving fluid is effected by means of functions which give the distribution of the fluid velocity $\vec{v} = \vec{v}(x, y, z, t)$ and of any two thermodynamic quantities pertaining to the fluid, for instance the pressure $p(x, y, z, t)$ and the density $\rho(x, y, z, t)$. All the thermodynamic quantities are determined by the values of any two of them, together with the equation of state; hence, if we are given five quantities, namely the three components of the velocity \vec{v} , the pressure p and the density ρ , the state of the moving fluid is completely determined.

All these quantities are, in general, functions of the coordinates x, y, z and of the time t . We emphasize that $\mathbf{v}(x, y, z, t)$ is the velocity of the fluid at a given point (x, y, z) in space and at a given time t , i.e., it refers to fixed points in space and not to specific particles of the fluid; in the course of time, the latter move about in space. The same remarks apply to ρ and p .

2.1 The velocity field

The velocity field $\vec{v}(\vec{x}, t)$ represents the center-of-mass velocity of the collection of molecules in a material particle with center-of-mass position \vec{x} at time t .

For a material particle of mass dM and volume dV , the total momentum is

$$d\vec{P} = \vec{v}(x, t) dM = \vec{v} \rho dV, \quad (1)$$

where $\rho(x, t)$ is the mass density field and $\vec{v}(x, t)$ is the center-of-mass velocity of the material particle. Because ρ and \vec{v} are averages over many molecules, their fluctuations become negligible provided the particle size is larger than the microscopic scale L'_{micro} .

2.1.1 Streamlines

Streamlines are curves everywhere tangent to the velocity field at a fixed time t_0 . They satisfy the ordinary differential equation

$$\frac{d\vec{x}}{dt} = \vec{v}(\vec{x}, t_0). \quad (2)$$

A streamline shows the instantaneous direction of flow and depends on the chosen reference frame. At a fixed time, there is one and only one streamline passing through each point in space, and streamlines cannot intersect.

2.1.2 Particle Trajectories

A particle trajectory (or particle orbit) is the actual path followed by an infinitesimal tracer particle advected by the flow. It satisfies

$$\frac{d\vec{x}}{dt} = \vec{v}(\vec{x}, t). \quad (3)$$

Given initial data $\vec{x}(t_0) = \vec{x}_0$, the trajectory is defined for all times. Unlike streamlines, different trajectories may intersect since they correspond to different times.

2.1.3 Streaklines

A streakline is formed by all particles that have passed through the same fixed point \vec{x}_0 , but at different emission times. Practically, this is obtained by continuously injecting dye or smoke at \vec{x}_0 .

For steady flow, where $\vec{v}(\vec{x}, t) = \vec{v}(\vec{x})$,

- streamlines,
- particle trajectories, and
- streaklines

all coincide.

For unsteady flow, these three types of flow visualization differ significantly, and streaklines may be misleading if interpreted as instantaneous flow directions.

2.2 Mass conservation

Let's consider some volume V_0 of space. The mass of fluid in this volume is $\int \rho dV$, where ρ is the fluid density, and the integration is taken over the volume V_0 . The mass of fluid flowing in unit time through an element $d\vec{S}$ of the surface bounding this volume is $\rho \vec{v} \cdot d\vec{S}$; the magnitude of the vector $d\vec{S}$ is equal to the area of the surface element, and its direction is along the normal. By convention, we take $d\vec{S}$ along the outward normal. Then $\rho \vec{v} \cdot d\vec{S}$ is positive if the fluid is flowing out of the volume, and negative if the flow is into the volume. The total mass of fluid flowing out of the volume V_0 in unit time is therefore

$$\oint_S \rho \vec{v} \cdot d\vec{S}, \quad (4)$$

where the integration is taken over the whole of the closed surface surrounding the volume in question.

Next, the decrease per unit time in the mass of fluid in the volume V_0 can be written

$$-\frac{\partial}{\partial t} \int \rho dV.$$

Equating the two expressions, we have

$$\frac{\partial}{\partial t} \int \rho dV = - \oint \rho \vec{v} \cdot d\vec{S}. \quad (5)$$

This is the global equation of mass conservation for an arbitrary fixed control volume.

The surface integral can be transformed by Green's formula to a volume integral:

$$\oint \rho \vec{v} \cdot d\vec{S} = \int \nabla \cdot (\rho \vec{v}) dV. \quad (6)$$

Thus

$$\int \left[\frac{\partial \rho}{\partial t} + \nabla \cdot (\rho \vec{v}) \right] dV = 0. \quad (7)$$

Since this equation must hold for any volume, the integrand must vanish, i.e.

$$\frac{\partial \rho}{\partial t} + \nabla \cdot (\rho \vec{v}) = 0. \quad (8)$$

This is the local equation of mass conservation, also called the *equation of continuity*. equation of continuity. Although derived from global mass conservation applied to fixed volumes, it is itself a local relation completely without reference to volumes. Expanding the expression $\nabla \cdot (\rho \vec{v})$, we can also write

$$\frac{\partial \rho}{\partial t} + \rho \nabla \cdot \vec{v} + \vec{v} \cdot \nabla \rho = 0. \quad (9)$$

The vector

$$\vec{j} = \rho \vec{v} \quad (10)$$

is called the *current density of mass* (or mass flux density). Its direction is that of the motion of the fluid, while its magnitude equals the mass of fluid flowing in unit time through unit area perpendicular to the velocity.

2.3 Incompressible flow

In a great many cases of the flow of liquids (and also of gases), their density may be supposed invariable, i.e. constant throughout the volume of the fluid and throughout its motion. In other words, there is no noticeable compression or expansion of the fluid in such cases. We then speak of incompressible flow.

Most liquids are perceived as incompressible under ordinary circumstances. A fluid is, however, effectively incompressible when flow speeds are much smaller than the velocity of sound.

All materials must nevertheless in principle be compressible. For in truly incompressible matter the sound velocity would be infinite, and that violates the relativistic injunction against any signal moving faster than the speed of light. Incompressibility is always an approximation and should, strictly speaking, be viewed as a condition on the flow rather than an absolute material property.

2.3.1 Global and local forms

Using $\rho(\vec{r}, t) = \rho_0 = \text{constant}$ in Eq. (5) we arrive at the global incompressibility condition

$$\oint_S \vec{v} \cdot d\vec{S} = 0, \quad (11)$$

Incompressible matter cannot accumulate anywhere, and equal volumes of incompressible fluid must enter and leave through any fixed surface per unit of time. Since the above condition refers to a single instant of time, it does in fact not matter whether the surface is fixed or moves in any way you may desire. On the other hand, using Eq. 8 we obtain the local incompressibility condition

$$\nabla \cdot \vec{v} = 0. \quad (12)$$

The divergence of the velocity field vanishes in incompressible flow. Notice that this local incompressibility condition does not refer to any volume of matter, only to the velocity field itself. A divergence-free field is sometimes called solenoidal.

2.3.2 Leonardo's Law.

Leonardo da Vinci observed that in an incompressible flow, the product of the cross-sectional area of a canal and the average velocity of the water remains constant. If A_1 and A_2 are two cross-sections of a canal and v_1, v_2 the corresponding average velocities, incompressibility implies

$$A_1 v_1 = A_2 v_2. \quad (13)$$

This follows from the fact that no fluid can pass through the canal walls. Applying the incompressibility condition to the closed surface composed of the two cross-sections and the canal walls gives

$$\oint \vec{v} \cdot d\vec{S} = \int_{A_2} \vec{v} \cdot d\vec{S} - \int_{A_1} \vec{v} \cdot d\vec{S} = 0.$$

The average velocity through any cross-section A is defined as

$$v = \frac{1}{A} \int_A \vec{v} \cdot d\vec{S}.$$

Hence, the quantity Av is constant along the canal for incompressible fluids, regardless of whether the flow is laminar or turbulent.

Leonardo's law, however, does *not* apply to compressible flows. In such cases the product of area and average velocity cannot remain constant, and a stronger law—accounting for density variations—is required.

2.4 Lagrangian and Eulerian derivatives

Fluid motion may be described in two complementary ways. In the *Lagrangian* viewpoint we follow a specific fluid element of mass Δm . Its motion is governed by Newton's law, and the time derivative d/dt represents the change experienced by that material element. This is the *material* or *Lagrangian* derivative. The velocity and acceleration of a fluid particle are written as

$$\vec{v} = \frac{D\vec{r}}{Dt}, \quad \vec{a} = \frac{D\vec{v}}{Dt}. \quad (14)$$

In the *Eulerian* viewpoint the fluid is treated as a field $f(\vec{r}, t)$ defined at fixed spatial positions. Changes in time may arise from variations at a fixed point, measured by $\partial f / \partial t$, or from the transport of fluid properties by the motion, represented by spatial derivatives.

According to differential calculus, a function $f(\vec{r}, t)$ has, in Cartesian coordinates, a differential of the form:

$$df = \frac{\partial f}{\partial t} dt + \frac{\partial f}{\partial x} dx + \frac{\partial f}{\partial y} dy + \frac{\partial f}{\partial z} dz = \frac{\partial f}{\partial t} dt + d\vec{r} \cdot \nabla f. \quad (15)$$

The temporal derivative of any fluid quantity is

$$\frac{df}{dt} = \frac{\partial f}{\partial t} + \vec{v} \cdot \nabla f. \quad (16)$$

Thus the relation between Eulerian and Lagrangian derivatives is

$$\frac{D}{Dt} = \frac{\partial}{\partial t} + \vec{v} \cdot \nabla. \quad (17)$$

Applied to the position vector,

$$\frac{D\vec{r}}{Dt} = \vec{v}, \quad (18)$$

showing that the material derivative of position yields the velocity. Applied to the velocity field,

$$\frac{D\vec{v}}{Dt} = \frac{\partial \vec{v}}{\partial t} + \vec{v} \cdot \nabla \vec{v}, \quad (19)$$

where the first term is the *local acceleration* and the second term is the *convective acceleration*, arising from velocity gradients in space.

For an incompressible fluid ($\nabla \cdot \vec{v} = 0$), the continuity equation gives

$$\frac{D\rho}{Dt} = 0 \quad \Rightarrow \quad \frac{\partial \rho}{\partial t} = -\vec{v} \cdot \nabla \rho.$$

This shows that although density may vary in space, each fluid element preserves its density as it moves.

2.5 Cauchy's equation

Applying Newton's Second Law to any comoving material particle with mass $dM = \rho dV$, its equation of motion becomes

$$dM \vec{a} = d\vec{\mathcal{F}}, \quad (20)$$

where $d\vec{\mathcal{F}}$ is the total force acting on the particle. Dividing by dV we arrive at the seductively simple equation, $\rho \vec{a} = \vec{f}^*$, where $\vec{f}^* = d\vec{\mathcal{F}}/dV$ is the *effective force density*. It was defined earlier:

$$\vec{f}_i^* = \vec{f}_i + \sum_j \nabla_j \sigma_{ij}, \quad (21)$$

where \vec{f}_i is the true body force density and σ_{ij} the stress tensor. Finally, inserting the acceleration field we have arrived at Cauchy's equation:

$$\rho \frac{D\vec{v}}{Dt} = \rho \left(\frac{\partial \vec{v}}{\partial t} + (\vec{v} \cdot \vec{\nabla}) \vec{v} \right) = \vec{f}^* \quad (22)$$

3 Incompressible frictionless (Euler) fluids

3.1 Euler equation for incompressible ideal flow

In an ideal homogeneous fluid, the only contact force is pressure, but in distinction to hydrostatics where the pressure is in balance with the body forces (like gravity), the effective density of force

$$\vec{f}^* = \vec{f}_e - \vec{\nabla}p \quad (23)$$

may now be non-vanishing. Inserted into the general dynamic equation we obtain, after dividing by the constant density $\rho(x, t) = \rho_0$,

$$\frac{\partial \vec{v}}{\partial t} + (\vec{v} \cdot \vec{\nabla})\vec{v} = \frac{\vec{f}_e}{\rho_0} - \frac{\vec{\nabla}p}{\rho_0}, \quad (24)$$

$$\vec{\nabla} \cdot \vec{v} = 0. \quad (25)$$

These two *Euler equations* govern the dynamics of incompressible ideal fluid, also called *Euler fluid*.

From the identity

$$\vec{\nabla}(\vec{A} \cdot \vec{B}) = (\vec{A} \cdot \vec{\nabla})\vec{B} + (\vec{B} \cdot \vec{\nabla})\vec{A} + \vec{A} \times (\vec{\nabla} \times \vec{B}) + \vec{B} \times (\vec{\nabla} \times \vec{A}), \quad (26)$$

with $\vec{A} = \vec{B} = \vec{v}$ it follows that

$$\frac{1}{2} \vec{\nabla}v^2 = \vec{v} \cdot \vec{\nabla}\vec{v} + \vec{v} \times (\vec{\nabla} \times \vec{v}), \quad (27)$$

so that, replacing $\vec{v} \cdot \vec{\nabla}\vec{v}$ in Euler's equation, we obtain

$$\rho_0 \frac{\partial \vec{v}}{\partial t} + \frac{\rho_0}{2} \vec{\nabla}v^2 - \rho_0 \vec{v} \times (\vec{\nabla} \times \vec{v}) + \vec{\nabla}p = \vec{f}_{\text{ext}}. \quad (28)$$

Unlike the hydrostatic case, it is no longer a necessary condition that external forces be derivable from a potential, i.e. that the forces be conservative. External forces may now include, in addition to electrical forces (including time-dependent ones), gravitational and inertial forces (centrifugal, Coriolis and linear acceleration), as well as magnetic forces.

In what follows we restrict ourselves to conservative forces, and more specifically to gravitational forces, centrifugal forces, and constant acceleration. Later, Coriolis forces will be introduced, although no conservative scalar potential exists for them. For the external forces we may then write

$$\vec{f} = -\rho \vec{\nabla}\Phi, \quad \text{with} \quad \Phi = \Phi_{\text{grav}} - \frac{\omega^2 r^2}{2} + \vec{a} \cdot \vec{r}.$$

Thus, the equation of motion for non-viscous fluids in the presence of these forces is

$$\rho_0 \frac{\partial \vec{v}}{\partial t} + \frac{\rho_0}{2} \vec{\nabla}v^2 - \rho_0 \vec{v} \times (\vec{\nabla} \times \vec{v}) + \vec{\nabla}p + \rho_0 \vec{\nabla}\Phi = 0. \quad (29)$$

3.2 Adiabatic and isentropic flow

The absence of heat exchange between different parts of the fluid (and also, of course, between the fluid and bodies adjoining it) means that the motion is adiabatic throughout the fluid. Thus the motion of an ideal fluid must necessarily be supposed adiabatic.

In adiabatic motion the entropy of any particle of fluid remains constant as that particle moves about in space. Denoting by s the entropy per unit mass, we can express the condition for adiabatic motion as

$$\frac{Ds}{Dt} = 0 = \frac{\partial s}{\partial t} + \vec{v} \cdot \vec{\nabla} s. \quad (30)$$

This is the general equation describing adiabatic motion of an ideal fluid. Using the continuity equation, we may write it as an “equation of continuity” for entropy:

$$\frac{\partial(\rho s)}{\partial t} + \vec{\nabla} \cdot (\rho s \vec{v}) = 0. \quad (31)$$

The product $\rho s \vec{v}$ is the *entropy flux density*.

The adiabatic equation usually takes a much simpler form. If, as usually happens, the entropy is constant throughout the volume of the fluid at some initial instant, it retains everywhere the same constant value at all times and for any subsequent motion of the fluid. In this case we can write the adiabatic equation simply as

$$s = \text{constant}. \quad (32)$$

Such a motion is said to be *isentropic*.

We may use the fact that the motion is isentropic to put the equation of motion in a somewhat different form. To do so, we employ the familiar thermodynamic relation

$$d\mathcal{H} = T ds + V dp, \quad (33)$$

where \mathcal{H} is the heat function per unit mass (enthalpy), $V = 1/\rho$ is the specific volume, and T is the temperature. Since $s = \text{constant}$, we have simply

$$d\mathcal{H} = V dp = \frac{dp}{\rho}. \quad (34)$$

Thus

$$\frac{\vec{\nabla} p}{\rho} = \vec{\nabla} \mathcal{H},$$

and Euler's equation can therefore be written in the form

$$\frac{\partial \vec{v}}{\partial t} + (\vec{v} \cdot \vec{\nabla}) \vec{v} = \frac{\vec{f}_e}{\rho_0} - \vec{\nabla} \mathcal{H}. \quad (35)$$

Or

$$\rho_0 \frac{\partial \vec{v}}{\partial t} + \frac{\rho_0}{2} \vec{\nabla} v^2 - \rho_0 \vec{v} \times (\vec{\nabla} \times \vec{v}) + \rho_0 \vec{\nabla} \mathcal{H} + \rho_0 \vec{\nabla} \Phi = 0, \quad (36)$$

$$\frac{\partial \vec{v}}{\partial t} + \vec{\nabla} \left(\frac{1}{2} v^2 + \mathcal{H} + \Phi \right) - \vec{v} \times (\vec{\nabla} \times \vec{v}) = 0. \quad (37)$$

From the isentropic condition $\vec{\nabla}\mathcal{H} = \vec{\nabla}p/\rho$ it follows, by taking the curl, that

$$\vec{\nabla}p \times \vec{\nabla}\rho = 0,$$

so that $p = p(\rho)$. That is, *isentropic flow is associated with a barotropic fluid*. Also, from

$$\rho \vec{\nabla}\mathcal{H} = \vec{\nabla}p,$$

taking the curl we obtain

$$\vec{\nabla}\rho \times \vec{\nabla}\mathcal{H} = 0,$$

from which $\mathcal{H} = \mathcal{H}(\rho)$.

The description of the state of a moving fluid is given by \vec{v} , p , and ρ , which altogether amount to five quantities. Five equations are therefore required, namely:

- Equation of motion (three equations when expressed in components),
- Continuity equation (one equation),
- Equation of state (one equation).

3.2.1 Noundary conditions

It is also necessary to prescribe boundary conditions. For an ideal fluid (non-viscous, without heat exchange) these are:

- (a) The fluid cannot penetrate a solid surface; that is, the component of the velocity perpendicular to the boundary is zero:

$$\vec{v} \cdot \hat{n}|_S = 0. \quad (38)$$

The vector \hat{n} is perpendicular to the surface. If the surface moves with velocity \vec{v}_S , then

$$\vec{v} \cdot \hat{n}|_S = v_S.$$

- (b) If there are two immiscible fluids, the pressure and the component of \vec{v} normal to the interface are the same on both sides. If the interface moves, then $\vec{v} \cdot \hat{n}|_S = v_S$ again.

3.2.2 irrotational flow

Now, taking the curl of the equation of motion, the pressure gradient term disappears and we obtain an equation involving only the velocity. It describes the time evolution of $\vec{\nabla} \times \vec{v}$, which we shall call the *vorticity*:

$$\frac{\partial}{\partial t}(\vec{\nabla} \times \vec{v}) - \vec{\nabla} \times [\vec{v} \times (\vec{\nabla} \times \vec{v})] = 0. \quad (39)$$

A solution—the simplest one—of this equation is

$$\vec{\nabla} \times \vec{v} = 0,$$

which is known as *irrotational flow*.

3.3 Bernoulli's equation

From the equation of motion, specialized to the steady state ($\partial \vec{v} / \partial t = 0$), taking the scalar product with $d\vec{r} = \vec{v} dt$ and noting that

$$d\vec{r} \cdot [\vec{v} \times (\vec{\nabla} \times \vec{v})] = \vec{v} \cdot [\vec{v} \times (\vec{\nabla} \times \vec{v})] dt = 0,$$

and since $\vec{v} \cdot (\vec{v} \times \vec{\xi}) = 0$, we conclude that

$$d\vec{r} \cdot \vec{\nabla} \left(\frac{1}{2} v^2 + \mathcal{H} + \Phi \right) = \frac{d}{dl} \left(\frac{1}{2} v^2 + \mathcal{H} + \Phi \right) = 0, \quad (40)$$

where d/dl denotes differentiation along the direction of the velocity, that is, along a streamline. It then follows that along a streamline the following condition holds, known as Bernoulli's equation:

$$H \equiv \frac{1}{2} v^2 + \mathcal{H} + \Phi = \text{constant}. \quad (41)$$

This field—which we shall call the Bernoulli field H —may be viewed as an extension of the effective potential to fluid in motion. The constant has the same value at all points on a given streamline, but may differ from one streamline to another. In the incompressible case, we have $\mathcal{H} = p/\rho$.

The tangent to a streamline at any point is in the direction of the velocity at that point. The equation of a streamline is therefore

$$\vec{v} \times d\vec{r} = 0.$$

In steady state these streamlines do not change with time and coincide with the trajectories of the fluid particles.

A *flow tube* is a curvilinear cylinder of moving fluid whose lateral walls are streamlines and whose end caps cut the lateral walls along curves made of fluid particles.

If a fluid of constant density flows in a horizontal tube, the term $\Phi = gz$ remains constant, so that Bernoulli's equation reduces to

$$\frac{1}{2} \rho v^2 + p = \text{constant}.$$

Consequently, in a horizontal tube, the greater the velocity, the lower the pressure, and vice versa. This effect is responsible for the lift of an airplane, since the shape of the wings is designed so that the curvature is greater on the upper surface. This causes the velocity over the wing to be higher and therefore the pressure lower. The resulting upward pressure force lifts the plane.

If A is the area of the wing and P_1 and P_2 are the pressures below and above the wing, respectively, the lift force is

$$F = A(P_1 - P_2) = \frac{A\rho}{2}(v_1^2 - v_2^2). \quad (42)$$

A useful approximation for the lift force is

$$F = \frac{A\rho}{2}(v_1^2 - v_2^2) = \frac{A\rho}{2}(v_1 + v_2)(v_1 - v_2) = A\rho \bar{v} (v_1 - v_2), \quad (43)$$

where the average velocity \bar{v} is very close to the speed of the airplane relative to the air.

Finally, note that the hydrostatic equation may also be obtained from Bernoulli's equation in the static case.

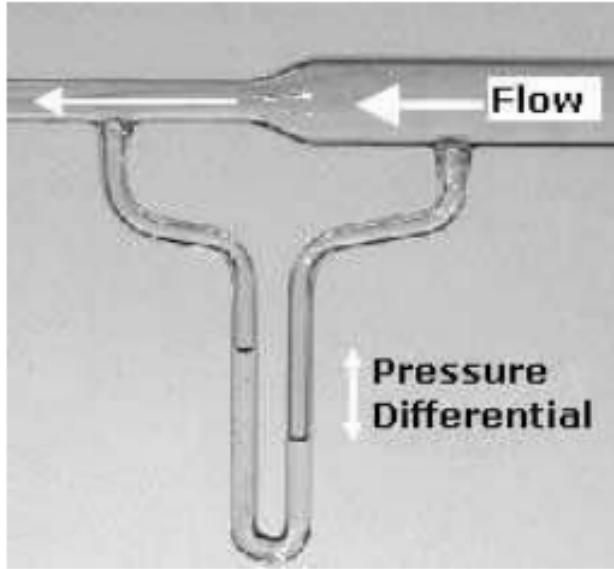


Figure 1: Venturi effect demonstration. Air streams from the right into a narrower tube where it speeds up. The shift in water level in the shunt tube underneath can be used to determine the pressure drop and thus the flow speed.

3.3.1 The Venturi effect

A simple duct with gently varying cross-section carries a constant volume flow rate Q of incompressible fluid (see Fig. 1). For simplicity we assume the duct is horizontal, such that gravity can be disregarded. For a streamline running through the duct, Bernoulli's theorem tells us that

$$H = \frac{1}{2}v^2 + \frac{p}{\rho_0} \quad (44)$$

takes the same value everywhere along the streamline. Approximating the velocity by its average value $U = Q/A$ over the cross-section A , the pressure becomes

$$p = \rho_0 \left(H - \frac{Q^2}{2A^2} \right). \quad (45)$$

This demonstrates the Venturi effect: the pressure rises when the duct cross-section increases, and conversely. The Venturi effect is used in many technological devices, from carburetors in car engines to perfume atomizers, as well as pumps and flow meters.

3.3.2 Torricelli's law

A barrel of wine has a little spout close to the bottom. If the plug in the spout is suddenly removed, the hydrostatic pressure accelerates the wine to stream out with considerable speed. Provided the spout is narrow compared to the size of the barrel, a nearly steady flow will soon establish itself.

Consider a streamline starting near the top of the barrel and running down through the middle of the spout. Near the top at height $z = h$, the fluid is almost at rest, $v \approx 0$. The pressure is

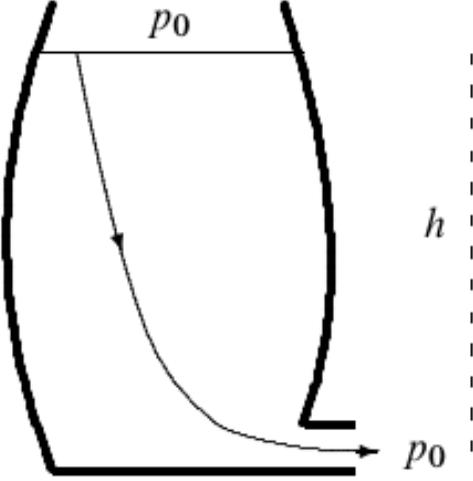


Figure 2: Wine running out of a barrel. The wine emerges with the same speed as it would have obtained by falling freely through the height h of the fluid in the barrel.

atmospheric, $p = p_0$, and the gravitational potential may be taken to be gz . Thus,

$$H_{\text{top}} = gh + \frac{p_0}{\rho_0}. \quad (46)$$

Just outside the spout, the fluid has horizontal velocity v and the pressure is also atmospheric. Hence,

$$H_{\text{bottom}} = \frac{1}{2}v^2 + \frac{p_0}{\rho_0}. \quad (47)$$

Equating the two expressions gives

$$\frac{1}{2}v^2 = gh, \quad (48)$$

with the solution

$$v = \sqrt{2gh}. \quad (49)$$

This result is called *Torricelli's law*. Surprisingly, the outflow speed equals the speed acquired by a freely falling particle over the same height.

3.3.3 Quasi-stationary emptying

If the barrel has constant cross-section A_0 (see Fig. 2), the average vertical flow velocity is $v_0 = vA/A_0$, where A is the cross-section of the spout and $v = \sqrt{2gz}$ the outflow velocity when the water level is z . Since $dz/dt = -v_0$, we obtain

$$\frac{dz}{dt} = -\frac{A}{A_0}\sqrt{2gz}. \quad (50)$$

Integration gives the emptying time

$$t_0 = \int_0^h \frac{dz}{\sqrt{2gz}} \frac{A_0}{A} = \frac{A_0}{A} \sqrt{\frac{2h}{g}}. \quad (51)$$

3.3.4 Case: The Pitot tube

The Pitot tube measures flow speed by converting kinetic pressure into static pressure. In steady state, water rises to height h so that

$$\frac{p_0}{\rho_0} = \frac{1}{2}U^2 + \frac{p}{\rho_0}, \quad (52)$$

which gives

$$\Delta p = \frac{1}{2}\rho_0 U^2 = \rho_0 gh. \quad (53)$$

Hence,

$$U = \sqrt{2gh}. \quad (54)$$

At a speed of 1 m s^{-1} the water rises 5 cm . Even if people do not know Bernoulli's theorem, farmers know better than to leave the barn door open in a storm: a strong gust reduces the pressure above the roof and can blow it off.

3.3.5 Rotating fluid in steady state

This problem has often been solved in a reference frame rotating with the liquid. Here we solve it in the inertial frame, by applying directly the equation of motion in steady state, with constant density ρ and in the presence of a uniform gravitational field.

In steady state it is true that

$$\vec{v} = \vec{\omega} \times \vec{r},$$

and therefore

$$\vec{\nabla} \times \vec{v} = 2\vec{\omega}, \quad \vec{v} \times (\vec{\nabla} \times \vec{v}) = 2(\vec{\omega} \times \vec{r}) \times \vec{\omega}.$$

A straightforward calculation then shows that the equation of motion reduces to

$$-\omega^2(x\hat{e}_x + y\hat{e}_y) + \frac{\vec{\nabla}p}{\rho} + g\hat{e}_z = 0. \quad (55)$$

From this, the three component equations follow:

$$\frac{\partial p}{\partial x} - \rho\omega^2x = 0, \quad (56)$$

$$\frac{\partial p}{\partial y} - \rho\omega^2y = 0, \quad (57)$$

$$\frac{\partial p}{\partial z} + \rho g = 0. \quad (58)$$

Consequently,

$$dp = \frac{\partial p}{\partial x}dx + \frac{\partial p}{\partial y}dy + \frac{\partial p}{\partial z}dz = \rho\omega^2(xdx + ydy) - \rho g dz, \quad (59)$$

from which

$$p = p_0 + \frac{\rho\omega^2}{2}(x^2 + y^2) + \rho g(z_0 - z), \quad (60)$$

in agreement with the result obtained in Hydrostatic chapter.

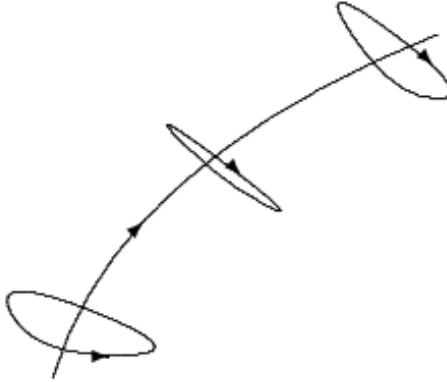


Figure 3: Around a vortex line there is local circulation of fluid.

3.4 Vorticity

The value of the Bernoulli field $H(x)$ at a point x depends only on the streamline passing through that point. Different streamlines will in general have different values of H . It is, however, often possible to relate the values of H for bundles of streamlines.

3.4.1 The vorticity field

Using the Euler equation for steady incompressible non-viscous flow, the gradient of the Bernoulli field including gravity becomes (in index notation)

$$\nabla_i H = \frac{1}{2} \nabla_i v^2 + \nabla_i \Phi + \frac{1}{\rho_0} \nabla_i p = v_j \nabla_i v_j - v_j \nabla_j v_i = (\vec{v} \times (\vec{\nabla} \times \vec{v}))_i. \quad (61)$$

Defining the *vorticity field* as the curl of the velocity field,

$$\vec{\omega} = \vec{\nabla} \times \vec{v}, \quad (62)$$

we obtain for steady flow

$$\vec{\nabla} H = \vec{v} \times \vec{\omega}. \quad (63)$$

The vorticity field goes back to Cauchy (1841) and is a quantitative measure of the local circulation in the fluid. In any region where the vorticity vanishes we have $\vec{\nabla} H = 0$, so the Bernoulli field must take the same value everywhere in that region, that is $H(x) = H_0$. A flow completely free of vorticity is said to be *irrotational*.

3.4.2 Vortex lines

The field lines of the vorticity field are called *vortex lines* (see Fig. 3), and are defined as curves that are everywhere tangent to the vorticity field. Like streamlines, they are solutions of the differential equation at a fixed time t_0 ,

$$\frac{d\vec{x}}{ds} = \vec{\omega}(\vec{x}, t_0), \quad (64)$$

where s is a running parameter along the curve. Unlike streamlines, the parameter is not time; here s has the dimensions of $(\text{length})^2/\text{time}$.

Like streamlines, vortex lines cannot cross each other, and since the vorticity field is a curl, it satisfies

$$\vec{\nabla} \cdot \vec{\omega} = 0.$$

Therefore, vortex lines can neither start nor end anywhere in the fluid; they must continue until they reach the boundary of the flow (possibly at infinity) or close upon themselves.

In steady flow, vortex lines are independent of the chosen time t_0 . Bernoulli's theorem, $(\vec{v} \cdot \vec{\nabla})H = 0$, follows immediately by dotting $\vec{\nabla}H = \vec{v} \times \vec{\omega}$ with \vec{v} . Similarly, dotting with $\vec{\omega}$ gives

$$(\vec{\omega} \cdot \vec{\nabla})H = 0,$$

showing that the Bernoulli field is also constant along vortex lines. Together these results show that the Bernoulli field is constant on the two-dimensional surfaces formed by combining vortex lines and streamlines. These are sometimes called *Lamb surfaces* or *Bernoulli surfaces*.

3.4.3 Equation of motion for vorticity

Since the vorticity field is derived from the velocity field, its equation of motion must follow from Euler's equation for the velocity field. Retracing the steps leading to $\vec{\nabla}H = \vec{v} \times \vec{\omega}$ but now allowing for time dependence, we obtain

$$\frac{\partial \vec{v}}{\partial t} = \vec{v} \times \vec{\omega} - \vec{\nabla}H. \quad (65)$$

This reduces to the steady case when $\partial \vec{v}/\partial t = 0$.

Taking the curl of both sides and using the identity that the curl of a gradient vanishes, we arrive at the vorticity equation

$$\frac{\partial \vec{\omega}}{\partial t} = \vec{\nabla} \times (\vec{v} \times \vec{\omega}). \quad (66)$$

A major lesson follows: if the vorticity vanishes initially, $\vec{\omega}(\vec{x}, t) = 0$ throughout a region V at time t , then $\partial \vec{\omega}/\partial t = 0$ there. Consequently, the vorticity field will not change, and if it is zero initially, it will remain zero forever. *Vorticity cannot be generated by the motion of an ideal fluid; it must be present from the outset.*

3.5 Circulation

Vorticity is a local property of a fluid. The corresponding global concept is called *circulation*, defined as the line integral of the velocity field along a closed curve C (see Fig. 4),

$$\Gamma(C, t) = \oint_C \vec{v}(\vec{x}, t) \cdot d\vec{\ell}. \quad (67)$$

If C encircles a whirling region, the circulation measures how strongly the fluid rotates around the loop. The sign depends on whether the circulation is taken with or against the whirling motion.

3.5.1 Stokes' theorem

The circulation of a vector field around a closed curve equals the flux of its curl through any surface S bounded by the curve:

$$\oint_C \vec{v} \cdot d\vec{\ell} = \int_S (\vec{\nabla} \times \vec{v}) \cdot d\vec{S}. \quad (68)$$

It does not matter which surface S is chosen, provided it has C as boundary and lies entirely within the fluid.

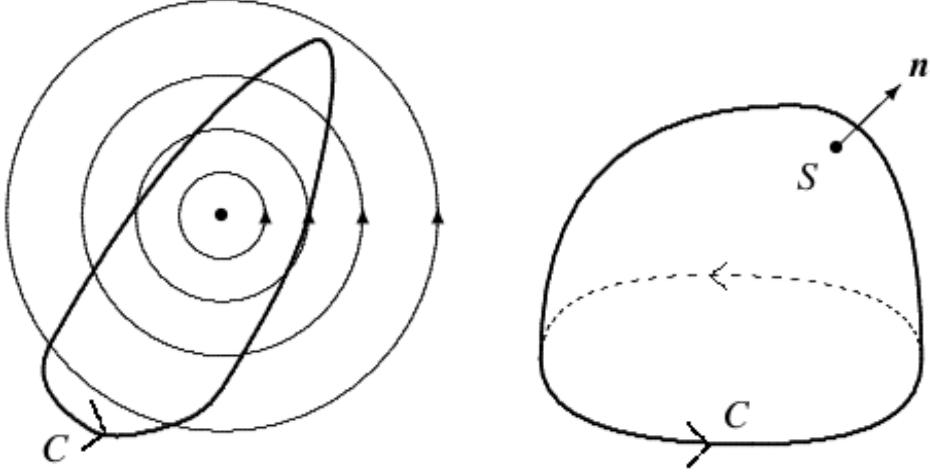


Figure 4: Left: A closed curve C encircling a whirl. Right: A surface S with perimeter C . The normal to the surface is consistent with the orientation of C (here using a right-hand rule).

3.6 Potential flow

In the absence of viscosity, a flow that is irrotational in a region will stay irrotational. High Reynolds numbers confine vorticity to thin boundary layers, so that the main body of the flow is nearly irrotational.

From $\vec{\nabla} \times \vec{v} = 0$ it follows that the velocity field is the gradient of a scalar potential,

$$\vec{v} = \vec{\nabla} \Psi. \quad (69)$$

The scalar Ψ is called the *velocity potential*.

3.6.1 Incompressible potential flow

In incompressible flow, $\vec{\nabla} \cdot \vec{v} = 0$ implies that Ψ satisfies Laplace's equation:

$$\nabla^2 \Psi = 0. \quad (70)$$

Inserting $\vec{\omega} = 0$ and $\vec{v} = \vec{\nabla} \Psi$ into Euler's equation gives

$$\vec{\nabla} \left(H + \frac{\partial \Psi}{\partial t} \right) = 0,$$

so the expression in parentheses depends only on time. Using the Bernoulli function and solving for the pressure, we obtain

$$p = C(t) - \rho_0 \left(\frac{1}{2} v^2 + \Phi + \frac{\partial \Psi}{\partial t} \right), \quad (71)$$

where $C(t)$ is an arbitrary function of time.

Potential flow in incompressible fluids is much simpler than flow with vorticity because Laplace's equation is well understood. Any linear superposition of solutions of Laplace's equation is again a solution, and practical problems reduce to finding the solution that satisfies the boundary conditions.

3.7 Application: Cylinder in uniform crosswind

A circular cylinder with axis along the z -axis and radius a is placed in an asymptotically uniform crosswind U along the x -axis. This does not break the translational invariance along the z -axis so that the velocity potential $\Psi(x, y)$ and the stream function $\psi(x, y)$ are both independent of z . Evidently, the natural choice is cylindrical coordinates (r, ϕ, z) .

3.7.1 Velocity potential

Asymptotically, for $r \rightarrow \infty$, the velocity potential must approach the field of a constant uniform crosswind, $\Psi \rightarrow Ux = Ur \cos \phi$. The linearity of Laplace's equation demands that the velocity potential is linear in the asymptotic values,

$$\Psi = U \cos \phi f(r), \quad (72)$$

where $f(r)$ is an unknown function that approaches r as $r \rightarrow \infty$.

Inserting this into the cylindrical Laplacian we have

$$\nabla^2 \Psi = \frac{\partial^2 \Psi}{\partial r^2} + \frac{1}{r^2} \frac{\partial^2 \Psi}{\partial \phi^2} + \frac{\partial^2 \Psi}{\partial z^2} + \frac{1}{r} \frac{\partial \Psi}{\partial r}, \quad (73)$$

which leads to

$$\frac{d^2 f}{dr^2} + \frac{1}{r} \frac{df}{dr} - \frac{f}{r^2} = 0. \quad (74)$$

Looking for power-law solutions $f \sim r^\alpha$, we find $\alpha = \pm 1$. Thus the general solution is

$$f(r) = Ar + \frac{B}{r},$$

with constants A and B .

The asymptotic condition implies $A = 1$, while B is determined by the boundary condition that the radial component of the velocity vanishes at the surface of the cylinder, i.e. $f'(a) = 0$. This yields $B = a^2$, and the solution becomes

$$\Psi = Ur \cos \phi \left(1 + \frac{a^2}{r^2} \right). \quad (75)$$

3.7.2 Velocity and pressure fields

The velocity field is obtained from derivatives of the potential. The non-vanishing components are

$$v_r = \frac{\partial \Psi}{\partial r} = U \cos \phi \left(1 - \frac{a^2}{r^2} \right), \quad (76)$$

$$v_\phi = \frac{1}{r} \frac{\partial \Psi}{\partial \phi} = -U \sin \phi \left(1 + \frac{a^2}{r^2} \right). \quad (77)$$

The radial velocity vanishes at the surface as expected, while the tangential component at $r = a$,

$$v_\phi|_{r=a} = -2U \sin \phi,$$

vanishes only at the stagnation points $\phi = 0, \pi$.

The pressure follows from Bernoulli's theorem. In the absence of gravity and normalized to vanish at infinity, it is

$$p = \frac{1}{2}\rho_0 (U^2 - v^2) = \frac{1}{2}\rho_0 U^2 \frac{a^2}{r^2} \left(4\cos^2 \phi - 2 - \frac{a^2}{r^2} \right). \quad (78)$$

On the surface $r = a$ this reduces to

$$p_a = \frac{1}{2}\rho_0 U^2 (4\cos^2 \phi - 3), \quad (79)$$

which is negative for $30^\circ < \phi < 150^\circ$.

The symmetry $p(\phi) = p(-\phi)$ implies that the total force in the y -direction (lift) vanishes, as expected. More surprisingly, the invariance $p(\pi - \phi) = p(\phi)$ implies that the total force in the x -direction (drag) also vanishes. This is known as *d'Alembert's paradox*.

3.7.3 Lift on a half-cylinder

The same solution also describes ideal flow over a cylinder that is half buried in a flat surface, since the normal velocity vanishes in the symmetry plane $y = 0$.

The vertical lift on a stretch of the half-cylinder of length L is

$$\mathcal{L}_y = - \int_0^\pi p_a dS_y = - \int_0^\pi p_a \sin \phi La d\phi = \frac{5}{3}\rho_0 U^2 La. \quad (80)$$

If the average density of the cylinder is $\rho_1 > \rho_0$, the lift-to-weight ratio is

$$\frac{\mathcal{L}_y}{Mg_0} = \frac{5}{3\pi} \frac{\rho_0}{\rho_1 - \rho_0} \frac{U^2}{ag_0}, \quad (81)$$

including buoyancy. There is therefore a critical wind speed beyond which the lift overcomes the weight and the half-buried cylinder can be lifted out of the ground.

4 Viscosity

4.1 Shear viscosity

Consider a fluid flowing steadily along the x -direction with a velocity field $v_x(y)$ that is independent of x but may vary with y . Such a field could, for example, be created by enclosing a fluid between moving plates, and is an elementary example of *laminar* or layered flow. If the velocity field has no y -dependence, so that the fluid is in uniform motion along the x -axis, there should not be any internal stresses. If, on the other hand, the velocity grows with y , so that $dv_x(y)/dy > 0$, we expect the fluid immediately above a plane $y = \text{const}$ to drag along the fluid immediately below because of friction and thus exert a positive shear stress, $\sigma_{xy}(y) > 0$, on this plane.

It seems reasonable to expect that a larger velocity gradient will evoke stronger stress. In Newton's law of viscosity the shear stress is made proportional to the gradient,

$$\sigma_{xy}(y) = \eta \frac{dv_x(y)}{dy}. \quad (82)$$

The proportionality constant η is called the coefficient of *shear viscosity*, the *dynamic viscosity*, or simply the *viscosity*. It is a measure of how strongly the moving layers of fluid are coupled by friction. In incompressible fluids there is also a bulk viscosity coefficient, but in practice it is usually much smaller and often negligible.

The viscosities of real fluids range over many orders of magnitude. Since dv_x/dy has dimension of inverse time, the unit for viscosity is Pa · s (pascal–seconds).

4.1.1 Molecular origin of viscosity in gases

In gases where molecules are far apart, internal stress arises from molecular bombardment of surfaces, transferring momentum across them. In liquids the stress is caused partly by molecular motion and partly by intermolecular forces. We therefore restrict the discussion to gases.

Molecules in a gas move nearly randomly at speeds much larger than the macroscopic flow velocity $v(x, t)$. In laminar planar flow with velocity $v_x(y)$ and gradient dv_x/dy , a molecule of mass m crossing a surface element dS_y from above carries more momentum in the x -direction than one crossing from below. The result is a net transfer of momentum downward.

Let λ be the mean free path and τ the typical time between collisions. Disregarding numerical factors of order unity, a layer of thickness λ above dS_y carries an excess momentum

$$dP_x \approx \rho \lambda^2 \frac{dv_x}{dy} dS_y. \quad (83)$$

The shear stress follows as momentum transferred per unit time and area:

$$\sigma_{xy} \approx \frac{dP_x}{\tau dS_y},$$

which recovers Newton's law with estimate

$$\eta \approx \rho \frac{\lambda^2}{\tau} \approx \rho \lambda v_{\text{mol}}. \quad (84)$$

Using $\lambda/\tau \approx v_{\text{mol}}$ and the rms molecular speed $v_{\text{mol}} = \sqrt{3RT}$, this explains the magnitude of gas viscosity.

4.1.2 Temperature dependence of viscosity

The viscosity of any material depends on temperature. Most liquids become thinner when heated (viscosity decreases), while gases become more viscous at higher temperatures.

For an ideal gas one finds

$$\eta = \eta_0 \sqrt{\frac{T}{T_0}}. \quad (85)$$

The viscosity of gases is independent of pressure and increases slightly faster than \sqrt{T} because of molecular attraction.

4.1.3 Kinematic viscosity

The estimate above suggests defining the *kinematic viscosity*

$$\nu = \frac{\eta}{\rho}. \quad (86)$$

It has units of $\text{m}^2 \text{s}^{-1}$. In ideal gases we have $\rho \sim p/T$, so that

$$\nu \sim \frac{T^{3/2}}{p},$$

and the kinematic viscosity decreases with temperature for isentropic flow.

Unlike the dynamic viscosity, it is the kinematic viscosity ν that appears in the Navier–Stokes equation. Under velocity-driven flows, air behaves as if it were more viscous than water by an order of magnitude, which explains why air seems “thicker” in certain aerodynamic contexts.

4.1.4 Velocity-driven planar flow

Assume laminar planar flow with velocity $v_x(y, t)$. Since there is no advection term $(\vec{v} \cdot \vec{\nabla})v_x$, Cauchy’s equation reduces to

$$\rho \frac{\partial v_x}{\partial t} = \frac{\partial \sigma_{xy}}{\partial y} = \eta \frac{\partial^2 v_x}{\partial y^2}. \quad (87)$$

Dividing by ρ gives

$$\frac{\partial v_x}{\partial t} = \nu \frac{\partial^2 v_x}{\partial y^2}. \quad (88)$$

This is a diffusion equation for velocity.

4.1.5 Case: Steady planar flow

In steady state,

$$\frac{\partial^2 v_x}{\partial y^2} = 0,$$

so

$$v_x = A + By.$$

For plates at $y = 0$ and $y = d$ with velocities 0 and U , respectively, boundary conditions yield

$$v_x(y) = \frac{U}{d}y. \quad (89)$$

The shear stress is

$$\sigma_{xy} = \eta \frac{U}{d}. \quad (90)$$

4.1.6 Case: Viscous friction

If the contact area is A and fluid thickness is d , the drag force is

$$D \approx \eta \frac{A}{d} U. \quad (91)$$

For a mass M experiencing only viscous drag,

$$M \frac{dU}{dt} = -\eta \frac{A}{d} U, \quad (92)$$

with solution

$$U = U_0 e^{-t/\tau}, \quad \tau = \frac{Md}{\eta A}. \quad (93)$$

The stopping distance is

$$L = \int_0^\infty U dt = U_0 \tau = \frac{U_0 M d}{\eta A}. \quad (94)$$

4.1.7 Case: Shear wave (Stokes' second problem)

Let a plate oscillate as $U(t) = U_0 \cos(\omega t)$. The velocity field is

$$v_x(y, t) = U_0 e^{-ky} \cos(\omega t - ky), \quad k = \sqrt{\frac{\omega}{2\nu}}. \quad (95)$$

This is a damped transverse wave. The wavelength is $\lambda = 2\pi/k = 2\sqrt{\pi\nu\tau}$ and the penetration depth $d = 1/k$.

A 1 kHz shear wave penetrates only 71 μm in air and 18 μm in water.

4.1.8 Case: Free momentum diffusion

The planar flow equation describes momentum diffusion. A Gaussian initial profile

$$v_x(y, 0) = U e^{-y^2/a^2}$$

evolves as

$$v_x(y, t) = U \frac{a}{\sqrt{a^2 + 4\nu t}} \exp\left(-\frac{y^2}{a^2 + 4\nu t}\right). \quad (96)$$

The width grows as $\sqrt{a^2 + 4\nu t}$. At late times,

$$\delta_{\text{front}} = 2\sqrt{\nu t}.$$

4.1.9 Case: Growth of a boundary layer (Stokes' first problem)

A plate at $y = 0$ starts moving at speed U . Seeking a similarity solution,

$$v_x(y, t) = U f\left(\frac{y}{\sqrt{\nu t}}\right), \quad (97)$$

leads to

$$f''(s) + \frac{1}{2} s f'(s) = 0. \quad (98)$$

The solution satisfying $f(0) = 1$, $f(\infty) = 0$ is

$$f(s) = \frac{1}{\sqrt{\pi}} \int_s^\infty e^{-u^2/4} du. \quad (99)$$

For large s , the decay is Gaussian,

$$v_x \sim \exp\left(-\frac{y^2}{4\nu t}\right),$$

as expected from momentum diffusion.

4.2 Dynamics of incompressible Newtonian fluids

Numerous everyday fluids obey Newton's law of viscosity (15.1), for example water, air, oil, alcohol, and antifreeze. A number of common fluids are only approximately Newtonian, for example paint and blood, and others are strongly non-Newtonian, for example tomato ketchup, jelly, and putty. There also exist *viscoelastic* materials that—depending on frequency—are both elastic and viscous. They are sometimes used in toys that can be slowly deformed like clay but also bounce like rubber balls when dropped on the floor.

In Newtonian fluids the shear stress σ_{xy} is directly proportional to the velocity gradient $\nabla_y v_x$ —also called the *shear strain rate*—with proportionality constant equal to the shear viscosity η . Most non-Newtonian fluids become “thinner” as the shear strain rate increases, meaning the shear stress grows slower than linearly. Even the most Newtonian of fluids, water, is thinner at shear strain rates above 10^{12} s^{-1} . Only a few fluids (for example some starches stirred in water) appear to thicken with increasing strain rate. The science of the general flow properties of materials is called *rheology*.

We shall in this section establish the general dynamics for incompressible, isotropic, and homogeneous Newtonian fluids, and postpone the analysis of the slightly more complicated compressible fluids to Section 15.5.

4.2.1 Isotropic viscous stress

Newton's law of viscosity (15.1) is a linear relation between the shear stress and the velocity gradient, only valid in a particular flow geometry. As for Hooke's law for elasticity, we want a definition of viscous stress that takes the same form in any flow geometry and in any Cartesian coordinate system.

Most fluids are not only Newtonian but also *isotropic*. In an isotropic fluid at rest there are no internal directions, and the stress tensor is determined by the pressure,

$$\sigma_{ij} = -p \delta_{ij}.$$

When such a fluid is set in motion, the velocity field $\vec{v}(\vec{x}, t)$ defines a direction at every point. Stress is produced by velocity gradients, hence viscous stresses are determined by the tensor $\nabla_i v_j$. In an incompressible fluid,

$$\sum_i \nabla_i v_i = \vec{\nabla} \cdot \vec{v} = 0,$$

so the most general symmetric tensor that can be built from velocity gradients is

$$\sigma_{ij} = -p \delta_{ij} + \eta (\nabla_i v_j + \nabla_j v_i). \quad (100)$$

This is the natural generalization of Newton's law of viscosity for incompressible flow in any Cartesian coordinate system. For steady planar flow $\vec{v} = (v_x(y), 0, 0)$ the only nonzero components are

$$\sigma_{xy} = \sigma_{yx} = \eta \nabla_y v_x,$$

confirming that η is indeed the shear viscosity.

4.3 The Navier–Stokes equations

The right-hand side of Cauchy's equation of motion equals the effective density of force

$$f_i^* = f_i + \sum_j \nabla_j \sigma_{ij}.$$

Using the stress tensor above and $\vec{\nabla} \cdot \vec{v} = 0$,

$$\sum_j \nabla_j \sigma_{ij} = -\nabla_i p + \eta \left(\sum_j \nabla_i \nabla_j v_j + \sum_j \nabla_j^2 v_i \right) = -\nabla_i p + \eta \nabla^2 v_i.$$

Assuming the fluid is homogeneous, i.e. η and ρ are constants, we obtain the Navier–Stokes equations for incompressible Newtonian fluids:

$$\frac{\partial \vec{v}}{\partial t} + (\vec{v} \cdot \vec{\nabla}) \vec{v} = \vec{g} - \frac{1}{\rho_0} \vec{\nabla} p + \nu \nabla^2 \vec{v}, \quad \vec{\nabla} \cdot \vec{v} = 0. \quad (101)$$

Here ρ_0 is the constant density, $\nu = \eta/\rho_0$ the kinematic viscosity, and $\vec{g} = \vec{f}/\rho_0$ the acceleration field of volume forces (e.g. gravity).

Although compact, these equations contain the full complexity of real fluid dynamics and generally cannot be solved in closed form except for highly symmetric geometries and idealized conditions.

4.4 Boundary conditions

The Navier–Stokes equation is first order in time and requires initial values and boundary conditions.

Density. The density may be discontinuous across a material interface and therefore provides no boundary condition.

Velocity. The normal component

$$v_n = \vec{v} \cdot \vec{n}$$

must be continuous across an interface between incompressible fluids. The tangential component

$$\vec{v}_t = \vec{n} \times (\vec{v} \times \vec{n})$$

must also be continuous, leading to the *no-slip condition*. The full velocity field is therefore continuous across any interface,

$$[\vec{v}] = 0. \quad (102)$$

An exception is an interface between a liquid and vacuum, where velocity may be discontinuous.

Stress. Newton's Third Law requires continuity of the stress vector:

$$[\boldsymbol{\sigma} \cdot \hat{n}] = 0. \quad (103)$$

At a vacuum boundary, $\boldsymbol{\sigma} \cdot \hat{n} = 0$, since vacuum cannot exert force.

Pressure. Pressure need not be continuous across all interfaces, but for incompressible viscous flow near a solid boundary it must coincide with the wall pressure.

Proof (sketch). Assume the wall lies locally at $z = Ax^2 + By^2 + 2Cxy$. Expanding the velocity near the wall,

$$\vec{v}(\vec{x}) = x\nabla_x \vec{v}_0 + y\nabla_y \vec{v}_0.$$

From $\vec{v} = 0$ at the wall we get all tangential derivatives zero at the origin. Using $\vec{\nabla} \cdot \vec{v} = 0$ we find $\nabla_z v_z = 0$, and from the stress tensor one obtains

$$\sigma_{zz} = -p.$$

Hence continuity of normal stress implies continuity of pressure at the wall.

4.5 Classification of flows

The most interesting phenomena in fluid dynamics arise from the competition between inertia and viscosity, represented in the Navier–Stokes equation by the advective acceleration $(\vec{v} \cdot \vec{\nabla})\vec{v}$ and the viscous diffusion term $\nu\nabla^2\vec{v}$. Inertia attempts to continue motion, whereas viscosity acts as a brake.

If inertia is dominant, the viscous term may be neglected and Euler’s equation for ideal flow is recovered. If viscosity dominates, the advective term may be dropped and the equations reduce to those describing slow *creeping flow*.

4.5.1 The Reynolds number

A dimensionless estimate of whether inertia or viscosity dominates is given by the Reynolds number, defined as the ratio of the advective to viscous terms. For a characteristic velocity scale U and length scale L ,

$$|\vec{\nabla}\vec{v}| \sim \frac{U}{L}, \quad |\nabla^2\vec{v}| \sim \frac{U}{L^2}.$$

Thus,

$$\text{Re} \approx \frac{|(\vec{v} \cdot \vec{\nabla})\vec{v}|}{|\nu\nabla^2\vec{v}|} = \frac{U^2/L}{\nu U/L^2} = \frac{UL}{\nu}. \quad (104)$$

Alternatively, the Reynolds number may be viewed as the ratio

$$\text{Re} \approx \frac{t_{\text{diff}}}{t_{\text{flow}}},$$

where $t_{\text{diff}} \sim L^2/\nu$ is the diffusion time and $t_{\text{flow}} \sim L/U$ is the flow time.

For $\text{Re} \ll 1$ the flow is slow and viscous (creeping). For $\text{Re} \gg 1$ inertia dominates and the flow is lively. At intermediate values, turbulence may develop.

4.5.2 Examples

Bathtub turbulence Typical values $U \sim 1$ m/s and $L \sim 1$ m give

$$\text{Re} \sim 10^6,$$

so turbulence is inevitable.

Curling For plates separated by $d \approx 43 \mu\text{m}$ and $U \approx 3 \text{ m/s}$,

$$\text{Re} = \frac{Ud}{\nu} \approx 140,$$

indicating laminar flow.

Water pipe For a pipe of diameter $d = 1.25 \text{ cm}$ and flow rate $Q = 100 \text{ cm}^3/\text{s}$,

$$U = \frac{Q}{\pi d^2} \approx 0.8 \text{ m/s}, \quad \text{Re} \approx 10^4,$$

which is deep in the turbulent regime.

System	Fluid	Size (m)	Velocity (m/s)	Reynolds number
Ship (Queen Mary 2)	Water	345	15	5.2×10^9
Submarine (Ohio class)	Water	170	12	2.0×10^9
Jet airplane (Boeing 747)	Air	71	250	1.2×10^9
Blue whale	Water	33	10	3.3×10^8
Car	Air	5	30	9.7×10^6
Swimming human	Water	2	1	2.0×10^6
Jogging human	Air	1	3	1.9×10^5
Herring	Water	0.3	1	3.0×10^5
Golf ball	Air	0.043	40	1.1×10^5
Ping-pong ball	Air	0.040	10	2.6×10^4
Fly	Air	0.01	1	6.5×10^2
Flea	Air	0.001	3	1.9×10^2
Gnat	Air	0.001	0.1	6.5
Bacterium	Water	10^{-6}	10^{-5}	10^{-5}

4.5.3 Hydrodynamic similarity

Two flows are said to be *hydrodynamically similar* if they have the same Reynolds number and congruent geometry.

For steady incompressible flow with $\vec{g} = 0$,

$$(\vec{v} \cdot \vec{\nabla})\vec{v} = -\frac{1}{\rho_0}\vec{\nabla}p + \nu\nabla^2\vec{v}. \quad (105)$$

Introduce dimensionless variables

$$\tilde{v} = U\tilde{v}, \quad \tilde{x} = L\tilde{x}, \quad p = \rho_0 U^2 \tilde{p}, \quad \tilde{\nabla} = \frac{1}{L}\tilde{\nabla}. \quad (106)$$

Then the Navier–Stokes equation becomes

$$(\tilde{v} \cdot \tilde{\nabla})\tilde{v} = -\tilde{\nabla}\tilde{p} + \frac{1}{\text{Re}}\tilde{\nabla}^2\tilde{v}. \quad (107)$$

The only remaining parameter is the Reynolds number. Thus, any two flows with equal Reynolds numbers and identical geometry have the same dimensionless solutions.

The viscous force on an object of size L scales as

$$D \sim \sigma L^2 \sim \eta \frac{U}{L} L^2 = \eta U L = \rho \nu U L = \rho U^2 L^2 \frac{1}{\text{Re}}.$$

For stones in air and water with equal Reynolds numbers,

$$\frac{D_{\text{air}}}{D_{\text{water}}} = \left(\frac{\eta \nu}{\eta \nu} \right)_{\text{air/water}} \approx 0.27.$$

4.5.4 Example: Robofly

Using hydrodynamic similarity, one can study insect flight using scaled models. A wing of size $L \approx 4$ mm flapping at 50 Hz has

$$U \approx \pi f L \approx 1.3 \text{ m/s}, \quad \text{Re} \approx 160.$$

The same flow can be reproduced using mineral oil of higher viscosity.

4.5.5 Example: High-pressure wind tunnels

In early flight research, Reynolds numbers were increased by raising pressure. Since η is pressure-independent and $\rho \propto p$,

$$\text{Re} \propto \rho.$$

The Variable Density Tunnel operated at 20 atm and enabled realistic testing on small-scale models.

Flows with the same Reynolds number may still differ in geometry. Hence Re predicts general behavior (laminar vs turbulent), not detailed flow structures.

References

- [1] A. Sepúlveda, *Hidrodinámica*. Fondo Editorial EIA, Colombia, 2019.
- [2] L. D. Landau and E. M. Lifshitz, *Fluid Mechanics*, vol. 6 of *Course of Theoretical Physics*. Pergamon Press, Oxford, UK, 2nd ed., 1987.
- [3] B. Lautrup, *Physics of Continuous Matter: Exotic and Everyday Phenomena in the Macroscopic World*. CRC Press, Boca Raton, 2 ed., 2011.