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HEAT OR MASS TRANSPORT TO POORLY STREAMLINED BODIES

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The heat and mass transfer between a stream and a streamlined body is examined taking into account the formation of turbulent fluid-filled zones behind the body.

As the stream velocity increases, the unseparated laminar flow around a body is replaced by separation, a vortex zone forms near the root domain of the body in which the mixing intensity increases as the Reynolds number grows. If the Re is not very large, the impurity heat or mass transfer to the body is determined entirely by the process of convective heat conductivity or diffusion, where the magnitude of the local flow on the surface decreases rapidly with distance from the inflow point. In this case the main transfer is realized in the frontal domain of the body and the root domain does not take part in the transfer in practice. As the Re increases, and as turbulence intensifies within the vortex zone, the role of this latter grows considerably and can become dominant for sufficiently large Re [1, 2].

The total heat for mass flux to the body is evidently comprised of flows in the laminar boundary layer domain up to its separation from the body surface and in the turbulized domain after separation. The transfer mechanisms differ substantially in the domains mentioned, and the construction of appropriate models for each requires the involvement of methods of different kinds. The transfer to the body frontal domain reached by the laminar boundary layer, and to the root domain adjoining the turbulized fluid is considered separately below. For definiteness, we shall speak about the stationary diffusion of impurity mass at constant concentrations far from the body and at its surface, but all the results to be obtained will be valid for heat transfer also.

Transfer to the Frontal Domain

The diffusion flow near the surface around which the laminar boundary layer flows can be found from the solution of the convective diffusion equation. If $Pe \gg 1$ was assumed, the known method of a thin diffusion boundary layer [1] is naturally used for the approximate solution. For $Sc \gg 1$ this layer is built-in into the hydrodynamic, for $Sc \ll 1$ on the other hand, the hydrodynamic boundary layer is built-in into the diffusion layer. The fluid velocity within the diffusion layer limits is evidently described by perfectly different relation-

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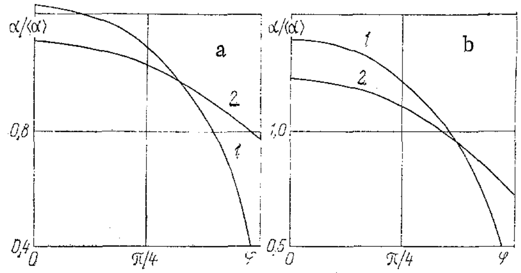

Fig. 1. Dependence of the relative local mass or heat transmission coefficient on φ at the frontal surface of a cylinder (a) and a sphere (b) for Sc $\geqslant$ 1 (1) and Sc $\ll$ 1 (2).

ships in these cases, and hence they should also be considered separately. As indicative examples, we consider below the transfer to a cylinder with axis oriented normally to the fluid flow, and to the sphere.

The Case Sc $\gg$ 1. Using the approximate solution of the boundary-layer equation for a cylinder, we write the approximate expressions for the tangential velocity component at the surface and for the stream function [3]:

$$v_\varphi \approx 2U\left(\frac{2UR}{\nu}\right)^{1/2} f\left(\frac{x}{R}\right)\frac{y}{R},\quad x = R\varphi,$$
$$\psi \approx UR\left(\frac{2UR}{\nu}\right)^{1/2} f\left(\frac{x}{R}\right)\left(\frac{y}{R}\right)^2,\quad f(\varphi) = (1.2326 - 0.4829\varphi^2)\varphi, \tag{1}$$

in which only the principal terms of the expansion in the coordinate y normal to the surface and the two first terms of the expansion in the coordinate x measured along the streamline at the surface from the stream inflow point $\varphi = 0$ are retained. Stream separation at $\varphi = \varphi_* \approx 91.5°$ corresponds to the approximation (1).

Considering just the self-similar solutions of the convective diffusion equation in the thin diffusion layer approximation, we obtain

$$c \approx \frac{c_0}{1.17}\int_0^z \exp\left(-\frac{4}{9}t^3\right)dt, \tag{2}$$

where we have introduced the self-similar variable

$$z = \frac{\mathrm{Re}^{1/2}\mathrm{Sc}^{1/3}}{2^{1/6}}\sqrt{f}\left(\int_0^\varphi \sqrt{f}\,d\varphi\right)^{-1/3}\frac{y}{R}. \tag{3}$$

The local diffusion stream to the surface is

$$j \approx 0.816\,(Dc_0/R)\,F(\varphi)\,\mathrm{Re}^{1/2}\mathrm{Sc}^{1/3}. \tag{4}$$

The function

$$F(\varphi) = [\varphi(1 - 0.3918\varphi^2)]^{1/2}\left(\int_0^\varphi [t(1 - 0.3918t^2)]^{1/2}dt\right)^{-1/3} \tag{5}$$

has been introduced in (3) and (4).

The total flux to the frontal part of the cylinder (up to separation) is

$$J_1 = 2R\int_0^{\varphi_*} j\,d\varphi \approx 2.40\,Dc_0\mathrm{Re}^{1/2}\mathrm{Sc}^{1/3} \tag{6}$$

and the corresponding Sherwood number is

$$\mathrm{Sh}_1 \approx 0{,}76\mathrm{Re}^{1/2}\mathrm{Sc}^{1/3}. \tag{7}$$

The dependence of the relative local mass elimination coefficient on the angle φ is presented in Fig. 1a.

For the flow around a sphere we have instead of (1) [3]

$$v_\varphi \approx \frac{3}{2} U\left(\frac{3UR}{\nu}\right)^{1/2} f\left(\frac{x}{R}\right)\frac{y}{R},\ x = R\varphi,$$
$$\psi \approx \frac{3}{4} UR^2\left(\frac{3UR}{\nu}\right)^{1/2} f\left(\frac{x}{R}\right)\left(\frac{y}{R}\right)^2 \sin\varphi, \tag{8}$$
$$f(\varphi) = (0{,}9277 - 0.3641\varphi^2)\varphi,$$

where $\varphi = \varphi_* \approx 91.5°$ again corresponds to flow separation.

The solution of the convective diffusion problem being considered is again expressed in the form of (2), but

$$z = \frac{\sqrt{3}}{2}\mathrm{Re}^{1/2}\,\mathrm{Sc}^{1/3}\,(f\sin\varphi)^{1/2}\left(\int_0^\varphi \sqrt{f}\sin^{3/2}\varphi d\varphi\right)^{-1/3}\frac{y}{R}. \tag{9}$$

In this case the local diffusion flux is represented thus

$$j \approx 0.731\,(Dc_0/R)\,F(\varphi)\,\mathrm{Re}^{1/2}\mathrm{Sc}^{1/3},$$
$$F(\varphi) = [\varphi(1 - 0{,}3925\varphi^2)\sin\varphi]^{1/2}\left(\int_0^\varphi [t(1 - 0{,}3925t^2)]^{1/2}\sin^{3/2}t dt\right)^{-1/3}, \tag{10}$$

while the total flux to the frontal surface of the sphere is

$$J_1 = 2\pi R^2 \int_0^{\varphi_*} j\sin\varphi d\varphi \approx 4{,}8Dc_0R\mathrm{Re}^{1/2}\mathrm{Sc}^{1/3}, \tag{11}$$

i.e., the corresponding Sherwood number is

$$\mathrm{Sh}_1 \approx 0{,}78\mathrm{Re}^{1/2}\mathrm{Sc}^{1/3}. \tag{12}$$

The dependence of $\alpha/\langle\alpha\rangle$ on φ is presented in Fig. 1b.

The Case $\mathrm{Sc} \ll 1$. In this case the diffusion boundary layer is considerably thicker than the hydrodynamic, i.e., we actually deal with convective diffusion to a body around which a potential ideal fluid stream flows. Problems of this kind (including nonstationary ones) were examined in [4-7]. For the local mass flux to a cylinder we have on the basis of [6, 7]

$$j \approx 0{,}798\frac{Dc_0}{R}\frac{\sin\varphi}{(1 - \cos\varphi)^{1/2}}\mathrm{Re}^{1/2}\mathrm{Sc}^{1/2}. \tag{13}$$

Correspondingly,

$$J_1 \approx 3.19Dc_0\mathrm{Re}^{1/2}\mathrm{Sc}^{1/2},\ \mathrm{Sh}_1 \approx 1.02\mathrm{Re}^{1/2}\mathrm{Sc}^{1/2}. \tag{14}$$

In place of (13) and (14), we have for a sphere [6, 7]

$$j \approx 1{,}197\frac{Dc_0}{R}\frac{\sin^2\varphi}{(2 - 3\cos\varphi + \cos^3\varphi)^{1/2}}, \tag{15}$$

$$J_1 \approx 7{,}1Dc_0R\mathrm{Re}^{1/2}\mathrm{Sc}^{1/2},\ \mathrm{Sh}_1 \approx 1.13\mathrm{Re}^{1/2}\mathrm{Sc}^{1/2}. \tag{16}$$

The dependence of $\alpha/\langle\alpha\rangle$ on φ for a cylinder and sphere are also presented in the figure for small Schmidt (or Prandtl) numbers.

Therefore, as the number Sc increases from values much less than one to values greatly exceeding it, the dependence of the diffusion fluxes on Re does not change but on Sc weakens: the appropriate exponent diminishes from 1/2 to 1/3.

Relationships of the form (13) and (15) had been used earlier to describe the transfer to the whole surface of a streamlined body in application to heat exchange with liquid metal fluxes [5] and to heat and mass transfer with filtration streams [6, 7]. In the second case these relationships reflect reality adequately, but this is not so in the first case. Indeed, for the validity of the relationships mentioned in the root domain it is necessary that the thermal boundary-layer thickness be much greater than the size of the vortex zone that is of the order R, which contradicts the thin layer approximation being used.

As is easy to see from Fig. 1, as the number Sc (Pr) diminishes the mass (heat) flux distribution on the frontal surface of a streamlined body becomes substantially more homogeneous, which is in agreement with the tendency noticed in [2], say.

Transfer in the Root Domain

Let us consider the vortex zone being formed behind the body to be filled with a fluid in a developed turbulent motion state, which reflects well the activity at high Reynolds numbers [3] when it is impossible to neglect transfer to the root part of the body surface more in comparison to the transfer to its frontal part [1, 2].

Since there is no sufficiently representative theory of turbulence as yet, it can be hoped to obtain a relationship, at this time, for the intensity of transfer which will be true just in order of magnitude. For this purpose it is necessary to use any of the existing semiempirical models of near-wall turbulence. Such models were developed mainly in application to conditions realizable in a turbulent boundary layer or in a turbulent flow in channels. In the vortex zone behind a body the hydrodynamic conditions are substantially different. Consequently, we use here model representations from [8], which are not associated with any assumptions about the nature of the flow at a distance from the surface. According to [8], the effective coefficient of turbulent kinematic viscosity in proximity to the wall is expressed in the form

$$\nu'(y) \approx \nu_0'(y)\Big(\int_0^{T(y)} R(t)dt \Big/ \int_0^{\infty} R(t)dt\Big)^2, \tag{17}$$

where $\nu_0'(y)$ is the turbulent kinematic viscosity at a distance y from the wall, determined in the usual way in terms of the effective values of the rms velocity fluctuations and mixing length realizable here, while T(y) is the characteristic "lifetime" of the vortices such that comparatively large vortices with $T > T(y)$ cannot approach the wall to a distance less than y.

Using the following approximation for the Lagrange correlation function

$$R(t) \approx \exp(-t/T_0(y)), \tag{18}$$

we obtain from (17)

$$\nu'(y) \approx \nu_0'(y)\,[1-\exp(-T(y)/T_0(y))]^2, \tag{19}$$

where $T_0(y)$ is the characteristic time scale for the vertices present at a distance y from the wall. Furthermore, it is assumed that $T(y)/T_0(y) = y/\Delta(y)$, where $\Delta(y)$ is the corresponding linear scale. It is clear that analogous relationships should hold even for the turbulent diffusion coefficients $D'(y)$ and $D_0'(y)$ and for the coefficients of turbulent thermal diffusivity.

Let us note that the already sufficiently rough approximation (18), that results in (19), permits obtaining automatically [8] both the known semiempirical model of Van Driest [9] for the viscous and transition layers of the turbulent boundary layer, and the Shablevskii model [10] for the turbulence miscibility domain.

The solution of the stationary diffusion equation under the evident assumption that the transverse component of the concentration gradient is significantly greater than the longitudinal, and with turbulent and molecular transport taken into account under the conditions c = 0 and y = 0, has the form

$$c \approx j\int_0^y \frac{dy}{D + D_0'[1-\exp(-y/\Delta)]^2} \approx j\int_0^y \frac{dy}{D + D'(y/\Delta)^2}. \tag{20}$$

The magnitude of the diffusion flow on the surface j, which can certainly depend on the coordinate x, is determined from the second boundary condition $c = c_0$ as $y \to \infty$. Therefore, the determination of j actually reduces to a problem in the representation of the quantities D_0' and Δ for vortices at different distances from the surface.

Let us introduce the characteristic thickness of the diffusion boundary layer in the root δ, and the minimal dimension λ of the turbulent vortices for which the order equality is valid (see [1], for example):

$$\lambda \sim \nu^{3/4} R^{1/4} U^{-3/4}. \tag{21}$$

The main resistance to the mass transfer process is evidently concentrated in the domain $y \leqslant \delta$.

Let $\delta \ll \lambda$. Then the diffusion boundary layer lies entirely within the hydrodynamic viscous sublayer, where the characteristic normal component of the velocity fluctuation has the order Uy/λ in conformity with the known hypothesis of L. D. Landau [1], and the mixing length is of the order y, i.e., $D_0' \sim Uy^2/\lambda$. It is clear that in this case $\Delta \sim \lambda$. Consequently, by solving (20) and requiring that the limit c(y) agree with c_0 as $y \to \infty$ we obtain with (21) taken into account

$$j \sim c_0 \left(\frac{U}{\lambda}\right)^{1/4} \frac{D^{3/4}}{\Delta^{1/2}} \sim \frac{Dc_0}{R} \mathrm{Re}^{13/16} \mathrm{Sc}^{1/4}. \tag{22}$$

Now, let $\delta \sim \lambda$. In this case, as before, $\Delta \sim \lambda$ but $D_0' \sim Uy$. After simple calculations there follows from (20) and (21):

$$j \sim c_0 U^{1/3} \frac{D^{2/3}}{\Delta^{2/3}} \sim \frac{Dc_0}{R} \mathrm{Re}^{5/6} \mathrm{Sc}^{1/3}. \tag{23}$$

If $\delta \gg \lambda$, but the outer boundary of the diffusion boundary layer is within the transition sublayer, then vortices that are greater than the minimal but are considerably smaller than the greatest, comparable in size to R, play the main role in mass transfer. In this case, as before, $D_0' \sim Uy$, but $\Delta \sim U\tau$, where $\tau \sim v^2/\varepsilon \sim (\nu R)^{1/2} U^{-3/2}$ is the characteristic vortex lifetime, which does not differ too radically from the minimal (here v is the characteristic velocity of such vortices, and ε is the specific dissipation of the turbulence energy). Then $\Delta \sim \lambda^{2/3} R^{1/3}$ and we obtain by the previous method from (20) and (21)

$$j \sim c_0 U^{1/3} \frac{D^{2/3}}{\Delta^{2/3}} \sim \frac{Dc_0}{R} \mathrm{Re}^{2/3} \mathrm{Sc}^{1/3}. \tag{24}$$

Finally, for $\delta \gg \lambda$, when the mass transfer is due to fluctuations of the large-scale vortices, $D_0' \sim UR$ and $\Delta \sim R$. In this case

$$j \sim c_0 (UR)^{1/2} \frac{D^{1/2}}{\Delta^{1/2}} \sim \frac{Dc_0}{R} \mathrm{Re}^{1/2} \mathrm{Sc}^{1/2}. \tag{25}$$

Let us note that it was actually assumed in the derivation of (22)-(25) that the exchange outside the diffusion layer is sufficiently intense so that the substance concentration at its outer boundary would be in agreement with its value in the undepleted flow.

The thickness δ of the diffusion layer can be estimated from the condition that the molecular and turbulent diffusion coefficients would be identical in order of magnitude for $y = \delta$. Hence

$$\delta/R \sim (\mathrm{Re} \cdot \mathrm{Sc})^{-1}. \tag{26}$$

The order inequality $\delta \leqslant \lambda$ here denotes $D \leqslant U\lambda$, i.e.,

$$\mathrm{Sc} \cdot \mathrm{Re}^{1/4} \geqslant 1, \tag{27}$$

while $\delta \geqslant \Delta \sim \lambda^{2/3} R^{1/3}$ is equivalent to $D \geqslant U\Delta$, which yields

$$\mathrm{Sc} \cdot \mathrm{Re}^{1/2} \leqslant 1. \tag{28}$$

Relationships (26)-(28) permit comprehension of the valid individual dependences (22)-(25), for which there are specific intervals of variation in the criteria Re and Sc, and for which their general form can be written:

$$J_2 = C_2 S_2 (Dc_0/R)\,\mathrm{Re}^m \mathrm{Sc}^n,\ \mathrm{Sh}_2 = C_2 \mathrm{Re}^m \mathrm{Sc}^n, \tag{29}$$

where C_2 is a factor on the order of one that depends on the shape of the streamlined body, S_2 is the area of the root part of its surface. As Re and Sc increase, the factor m grows from 1/2 to a value equal to approximately 0.8, while the factor n drops from 1/2 to 1/4. If $\mathrm{Sc} \ll 1$, $\mathrm{Re}\cdot\mathrm{Sc} \sim 1$, then the dependence of Sh_2 on Re and Sc is the same in its structure as is the analogous dependence for Sh_1.

The relationships (7), (12), (14), and (16) from the frontal domain of the body and (29) for the root permit construction of semiempirical formulas of the form

$$\mathrm{Sh} = 1 + s_1 C_1 \mathrm{Re}^{1/2} \mathrm{Sc}^k + s_2 C_2 \mathrm{Re}^m \mathrm{Sc}^n \tag{30}$$

for the Sherwood number referred to the whole body. The coefficients C_1, k, m, and n for bodies of simple shape can here be considered known, while the coefficient C_2 should be determined experimentally. During such a determination of C_2 only data of experiments conducted for Re and Sc from ranges corresponding to the m and n values under consideration should be used. This requirement was not satisfied, for instance, in obtaining the empirical formula of B. D. Katsnel'son and F. A. Timofeeva for the heat transfer of a flow to a spherical body when test data were used that referred to the most diverse values of Sc (Pr) [2].

With the exception of the case $\mathrm{Re}\cdot\mathrm{Sc} \sim 1$, the exponent $m > 1/2$, i.e., for sufficiently large Re the role of the root domain in mass or heat transfer actually becomes governing, as much experimental data indicate [1, 2].

Results obtained above for $\mathrm{Sc} \leqslant 1$ are confirmed fairly by data on the heat transfer between a body and a stream including the mentioned empirical Katsnel'son–Timofeeva formula. For $\mathrm{Sc} \gg 1$ the theory of transfer to a solid surface from a turbulized fluid is confirmed in qualitative respects by the I. A. Bagotskaya tests, for example, on the diffusion to a rotating disc electrode in electrochemical reactions of oxygen reduction and hydrogen evolution, as well as by the experiments of A. I. Fedorova and G. L. Vidovich on the diffusion to a moving flat plate for the reaction of cathodic reduction of iodine, described in detail in [1]. In particular, a completely single-valued deduction was made about the increase in the exponent as a function of Sh and Re as Re increased from 0.5 to a value on the order of 0.8 or even higher.

NOTATION

C_i, coefficients in (30); c, concentration; c_0, value of c at a distance from the body; D, molecular diffusion coefficient; D_0', D', turbulent diffusion coefficients without and with taking account of the influence of the solid surface; F, f, functions determined in (5), (10) and in (2), (8); J, j, total and local diffusion flows; k, m, n, exponents in (29) and (30); R, cylinder or sphere radius; R(t), Lagrange correlation function; S, body surface area; s, fraction of the area; T_0, T, characteristic time scales; τ, fine vortex time scale; U, flow velocity; v_φ, tangential velocity component near the body surface; x, y, tangential and normal coordinates; α, coefficient of heat or mass transfer; Δ, characteristic linear turbulent scale; δ, diffusion layer thickness; λ, minimal vortex size; ν, molecular kinematic viscosity; ν_0', ν', turbulent kinematic viscosity without and with the influence of the surface taken into account; φ, the angular coordinate; ψ, stream function; $\mathrm{Re} = UR/\nu$; $\mathrm{Sc} = \nu/D$; $\mathrm{Pe} = \mathrm{Re}\cdot\mathrm{Sc}$; $\mathrm{Sh} = JR/Dc_0S$; the subscripts 1 and 2 refer to the frontal and root parts of the body surface.

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PRECIPITATION OF A CLOUD OF HEATED PARTICLES ON A HORIZONTAL PLANE

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The precipitation of a cloud of heated monodisperse particles in a field of external force is numerically investigated.

The nature of the motion of a cloud of particles in an infinite medium under the action of an external force (gravity) depends on the degree of hydrodynamic interaction between the particles, which is mediated by the carrier phase [1-3]. If the particle concentration in the cloud is low, each particle moves like a solitary particle, independently of the others (the "filtration" regime). If the particle concentration is sufficiently high this regime gives way to the "entrainment" regime, where the medium between the particles is involved in the motion; the assembly of particles moves at a speed that exceeds the speed of fall of a single particle.

In [4, 5] the motion and precipitation of a cloud possessing plane symmetry (in a direction perpendicular to the action of the external force one of the dimensions of the cloud was much greater than the other) were examined. A numerical solution of the unsteady two-dimensional equations of motion of a two-phase medium showed that the interaction of the cloud with the precipitation surface depends significantly on the regime of motion. In the filtration regime all the particles move in straight trajectories, perpendicular to the precipitation plane, and fall on this plane in the region of initial projection of the cloud; the final distribution of the precipitated particles does not depend on the initial height of the cloud. Entrainment motion of the cloud produces a large-scale vortical flow of the carrier medium in the form of two symmetric cylindrical eddies, which increase in size (the analog of this solution in the case of axial symmetry is a vortical ring — a torus). In the plane of symmetry the gas moves downward in the direction of the external force and on the periphery it rises. The particles are involved in this motion and become concentrated in the cores of the eddies; thus, most of the particles move in directions perpendicular to that of the external force. This effect becomes more pronounced as the cloud approaches the precipitation plane owing to the spreading of the gas along it, which causes additional sideways transport of particles. This leads to precipitation of some particles at great distances from the plane of symmetry, exceeding the initial radius of the cloud.

In [4, 5] the precipitation process was considered in the isothermal case, where the gas and particle temperatures are equal. In this paper we examine, in a plane formulation, the precipitation of a cloud of particles in the nonisothermal case, where the initial particle temperature exceeds the temperature of the surrounding gas. At the initial time a cold gas, in static equilibrium in a field of external force, contains a motionless cloud of heated solid or liquid spherical particles (a monodisperse aerosol). In describing the motion of the disperse medium we adopt the usual assumptions of mechanics of heterogeneous media [6], and regard the gas and particles as two interpenetrating and interacting continua. We consider systems in which the volume fraction of the particles and the ratio of the true gas and particle densities are small; collisions, breakup, and evaporation of the particles are insignificant; viscous dissipation of energy is negligible; the temperature is assumed constant through-

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