

Article

A Phenomenological Interpretation of Laminar Natural Convection Correlations

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Abstract

Natural convection heat transfer is commonly described using dimensionless correlations relating the Nusselt number to the Grashof and Prandtl numbers. For gases and non-metallic liquids, these parameters combine into the Rayleigh number, $Ra = Gr \cdot Pr$, while for liquid metals they form the Boussinesq number, $Bo = Gr \cdot Pr^2$, whose physical roles are often underrated. Focusing on the classical configuration of laminar flow over an isothermal vertical plate, the present work adopts a phenomenological perspective to examine how the governing variables influence the thickness of the thermal boundary layer, and, consequently, the heat transfer rate. Within this framework, the Rayleigh and Boussinesq numbers emerge naturally as the controlling dimensionless parameters. Once a power law functional dependence between the Nusselt number and the Rayleigh and Boussinesq numbers is assumed, a physically consistent range for the associated exponents is identified, together with the necessity of introducing a weak corrective function of the Prandtl number, which plays just a refining role. An overall order-of-magnitude analysis is further developed to recover the classical one-fourth power law structure of natural convection correlations and to clarify the origin of the Prandtl number corrective function. The primary contribution of this study is to demonstrate that the Rayleigh and Boussinesq numbers are not merely the outcome of solving the governing equations, but are the dimensionless groups that encapsulate the underlying physics of natural convection, in contrast to the Grashof and Prandtl numbers considered separately. This perspective provides a comprehensive physical interpretation of the existing heat transfer correlations, also offering guidance for identifying the appropriate dimensionless parameters when developing new correlations, useful both to researchers and educators.

Keywords: natural convection; Rayleigh number; Boussinesq number; Grashof number; Prandtl number; thermal boundary layer; phenomenological approach



Academic Editor: Maria Vicidomini

Received: 20 January 2026

Revised: 16 February 2026

Accepted: 18 February 2026

Published: 21 February 2026

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1. Introduction

Natural convection heat transfer originates from fluid motion induced by buoyancy forces generated by temperature differences, in the absence of any externally imposed flow. A canonical configuration used to investigate this phenomenon is the laminar flow developing along an isothermal vertical flat plate maintained at a uniform temperature T_p higher than that of the surrounding quiescent fluid, T_∞ . This geometry provides a physically transparent framework for examining the interplay between thermal and hydrodynamic boundary layers.

Despite the long history of research on natural convection adjacent to vertical surfaces, the subject continues to receive considerable attention owing to its relevance in modern technological applications. These include passive thermal management of electronic devices, energy-efficient building systems, solar thermal collectors, safety-related thermal processes in advanced nuclear systems, and aerospace applications, as discussed, for example, in [1–5]. In parallel, renewed interest has emerged in the fundamental theoretical aspects of natural convection, as documented in the recent comprehensive review authored by Fan et al. [6].

Although the isothermal vertical plate is adopted here as a reference configuration, the physical interpretations derived from this model extend to a broader class of buoyancy-driven flows, since they reflect general features of natural convection boundary layer development under the assumptions of laminar flow, Boussinesq approximation applicability and constant physical properties. Indeed, as discussed by Churchill [7], laminar heat transfer correlations for vertical plates can often be extended to different geometrical configurations and thermal boundary conditions by appropriately redefining both the characteristic length and the reference temperature.

For laminar flow, the typical approach to derive correlations for natural convection from a vertical plate—readily available in nearly all textbooks; see, for example, Gebhart [8] and Incropera et al. [9]—is based on the numerical solutions of the boundary layer equations. In particular, Ostrach [10] solved the coupled momentum and energy equations for Prandtl numbers ranging from 0.01 to 1000, a range representative of mainly gases and non-metallic liquids. Later, Sparrow and Gregg [11] extended these solutions to lower Prandtl numbers, thereby addressing the behavior of liquid metals.

From these solutions, the wall temperature gradient and the associated heat flux can be evaluated, leading to the well-known relationship

$$Nu = F(Pr) Gr^{1/4} \quad (1)$$

where the function $F(Pr)$ accounts for the influence of the Prandtl number. Various analytical expressions for $F(Pr)$ have been proposed by interpolating numerical solutions or by constructing empirical fitting formulas. Among these, the correlation proposed by Ede [12],

$$F(Pr) = \frac{0.6 Pr^{1/2}}{(1 + 2Pr^{1/2} + 2Pr)^{1/4}} \quad (2)$$

and that by Churchill and Usagi [13],

$$F(Pr) = \frac{0.503 Pr^{1/4}}{\left[1 + \left(\frac{0.492}{Pr}\right)^{9/16}\right]^{4/9}} \quad (3)$$

are the most widely adopted.

In the asymptotic limit $Pr \rightarrow \infty$, both Equations (2) and (3) reduce to a power-law dependence of the form $F(Pr) = C_1 Pr^{1/4}$, consistent with the analytical result obtained by LeFevre [14]:

$$Nu = 0.5027 (Gr Pr)^{1/4} \quad (Pr \rightarrow \infty) \quad (4)$$

Conversely, for moderate-to-high Prandtl numbers, the function $F(Pr)$ may be written as

$$F(Pr) = f_1(Pr) Pr^{1/4} \quad (5)$$

where $f_1(Pr)$ is a weak function of the Prandtl number. As illustrated in Figure 1, this function varies only moderately over a wide range of Pr , which justifies the use of the constant average value of 0.4425 for engineering estimates, as suggested by McAdams [15].

Combining Equations (1) and (5), the heat transfer correlation for gases and non-metallic liquids can be expressed as

$$Nu = f_1(Pr) (Gr Pr)^{1/4} = f_1(Pr) Ra^{1/4} \quad (Pr \rightarrow \infty) \quad (6)$$

where the Rayleigh number is defined as

$$Ra = Gr Pr = \frac{g\beta(T_p - T_\infty)x^3}{\alpha\nu} \quad (7)$$

This formulation highlights that, for moderate-to-high Prandtl numbers, the Rayleigh number emerges as the dominant dimensionless parameter governing natural convection.

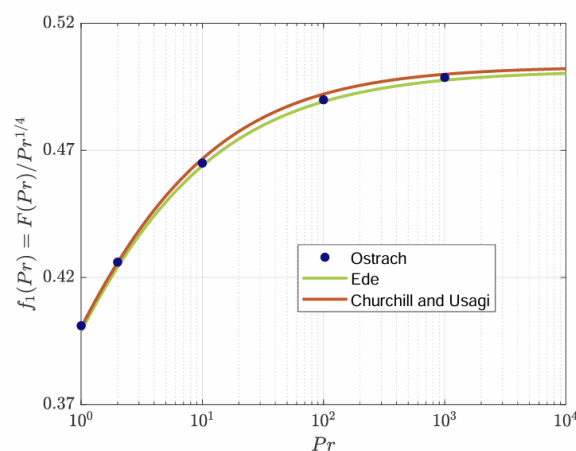


Figure 1. Function $f_1(Pr)$ according to the fitting formulas proposed by Ede [12], and by Churchill and Usagi [13], together with the numerical similarity solutions provided by Ostrach [10].

A different asymptotic behavior is observed in the limit $Pr \rightarrow 0$. In this case, Equations (2) and (3) approach $F(Pr) = C_2 Pr^{1/2}$, leading to the asymptotic result reported by LeFevre [14]:

$$Nu = 0.6004 Gr^{1/4} Pr^{1/2} = 0.6004 (Gr Pr^2)^{1/4} \quad (Pr \rightarrow 0) \quad (8)$$

For low but finite Prandtl numbers, the function $F(Pr)$ can be conveniently rewritten as

$$F(Pr) = f_2(Pr) Pr^{1/2} \quad (9)$$

where $f_2(Pr)$ is again a weak function of Pr , as shown in Figure 2. Following Sparrow and Gregg [11], $f_2(Pr)$ can be reasonably approximated by the constant average value of 0.565 for many practical applications. The corresponding heat transfer correlation then becomes

$$Nu = f_2(Pr) (Gr Pr^2)^{1/4} = f_2(Pr) Bo^{1/4} \quad (Pr \rightarrow 0) \quad (10)$$

where the Boussinesq number is defined as

$$Bo = Gr Pr^2 = \frac{g\beta(T_p - T_\infty)x^3}{\alpha^2} \quad (11)$$

This representation emphasizes that, for liquid metals, the Boussinesq number plays a role analogous to that of the Rayleigh number for gases and non-metallic liquids.

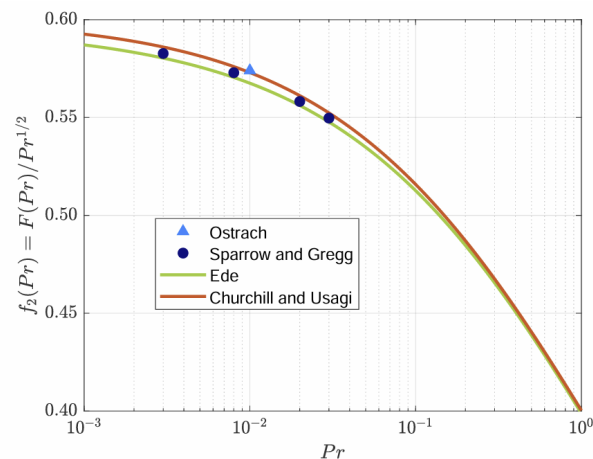


Figure 2. Function $f_2(Pr)$ according to the fitting formulas proposed by Ede [12], and by Churchill and Usagi [13], together with the numerical similarity solutions provided by Ostrach [10], and by Sparrow and Gregg [11].

The central importance of the Rayleigh and Boussinesq numbers was first systematically emphasized by Bejan [16], who employed a scale analysis of the coupled boundary layer equations. Subsequent studies extended this framework to fluids with intermediate Prandtl numbers and to unsteady flow conditions, recovering the steady-state as a limiting case, as shown, for example, by Lin and co-workers [17,18], who followed the approach pioneered by Patterson and Imberger [19]. Additional studies dealing with scale analyses applied to natural convection from vertical surfaces subjected to different boundary conditions have been executed by Capobianchi and Aziz [20], and Lin et al. [21]. Recent direct numerical simulations further confirmed the $Nu \sim Ra^{1/4}$ and $Nu \sim Bo^{1/4}$ relationships—see, e.g., Shishkina [22], Wang et al. [23], and Chand [24].

Within this general context, it is evident that the available theoretical treatments of laminar natural convection correlations rely on either similarity solutions or scale analyses of the governing differential equations. While these approaches have provided rigorous and widely accepted results, a purely phenomenological reconstruction of the correlations appears to be lacking. Therefore, the objective of the present study is to provide a phenomenological interpretation of the laminar natural convection correlations expressed by Equations (6) and (10). The analysis builds upon a phenomenological boundary layer framework previously developed for forced convection by our research group [25], and seeks to elucidate how each independent variable qualitatively affects the thickness of the thermal boundary layer, which constitutes the thermal resistance to the heat flux. By combining these qualitative effects, the Rayleigh and Boussinesq numbers—rather than the Grashof and Prandtl numbers separately—naturally arise as the most appropriate dimensionless groups to describe the phenomenology of natural convection. Of course, as our approach does not rely on quantitative equations, it cannot specify a precise functional dependence between the Nusselt number and the Rayleigh and Boussinesq numbers. On the other hand, once the power law dependence is adopted, physically consistent bounds for the associated power law exponents and a clear interpretation of the weak corrective functions of the Prandtl number are provided.

Overall, the novelty of the present work is not to provide new results, but rather to give a phenomenological interpretation of widely adopted heat transfer correlations, offering a different perspective on the topic and a new pedagogical approach useful to discuss the structure of natural convection heat transfer correlations without requiring any advanced mathematical knowledge.

Two appendices complement the main analysis by presenting an overall order-of-magnitude treatment that reproduces the one-fourth power-law dependence of the Nusselt number on the Rayleigh and Boussinesq numbers and further clarifies the origin of the Prandtl-number corrective functions.

2. The Phenomenology Behind the Heat Transfer Correlations for the Isothermal Vertical Plate

An isothermal vertical flat plate, maintained at a temperature T_p higher than the temperature T_∞ of the undisturbed fluid, is considered. Under the no-slip condition, the fluid layer adhering to the plate surface is at the plate temperature T_p , and heat is transferred through successive adjacent fluid layers in relative motion. As a consequence, a temperature profile develops in the direction normal to the plate, decreasing monotonically from T_p at the surface and asymptotically approaching the undisturbed fluid temperature T_∞ . The distance from the plate at which the temperature difference with respect to T_∞ becomes negligible defines the conventional thickness of the thermal boundary layer, denoted as δ_t . In the following, the effect of each independent variable on δ_t is examined in order to determine the convection heat transfer coefficient, first for moderate-to-high- Pr fluids (gases and non-metallic liquids) and subsequently for low- Pr fluids (liquid metals).

2.1. The Nu - Ra Correlation for Gases and Non-Metallic Liquids

Let x denote the vertical coordinate aligned with the upward bulk motion parallel to the plate, with $x = 0$ defined at the leading edge, and let z denote the coordinate normal to the plate, with $z = 0$ defined at the plate surface.

As the fluid moves upward, its temperature progressively increases due to the heat received by the plate. This leads to an enlargement of the region affected by temperature gradients and a corresponding “stretching” of the temperature profile in the direction normal to the plate. Therefore, the thickness δ_t of the thermal boundary layer, which originates at the leading edge of the plate, increases with increasing x . On the other hand, since the temperature difference between the plate and the adjacent fluid diminishes with increasing x , the reduced heat transfer rate results in a progressively less pronounced stretching of the temperature profile, leading to a decrease in $d\delta_t/dx$. A sketch of the temperature profile and the corresponding thermal boundary layer is shown in Figure 3.

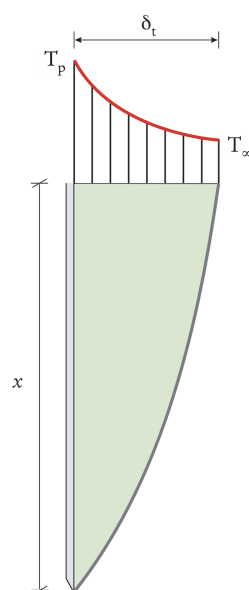


Figure 3. The temperature profile and the thermal boundary layer for the heated vertical plate.

The thickness of the thermal boundary layer also increases as the fluid's thermal diffusivity $\alpha = k/(c_p\rho)$ increases. Indeed, as α increases, temperature variations occur more rapidly, since heat transfer effects, proportional to the fluid's thermal conductivity k , become predominant over thermal inertia, quantified by $c_p\rho$. Therefore, over the time interval required to cover a given distance x , the fluid temperature increases with α , causing the temperature profile to stretch farther from the plate, which results in an increased thickness of the thermal boundary layer.

Conversely, δ_t increases as the average fluid velocity within the thermal boundary layer decreases. This is because a lower average velocity results in a longer traveling time, which allows for a greater temperature increase. Consequently, a more significant stretching of the temperature profile in the direction normal to the plate and a related thickening of the thermal boundary layer occur. In this regard, since the average fluid velocity within the thermal boundary layer is not known a priori, its value can be assumed to be directly related to a velocity U_c characteristic of natural convection. This characteristic velocity can be estimated by considering a unit mass of fluid at the average temperature T_{av} , surrounded by undisturbed fluid at temperature T_∞ and ideally subjected solely to the buoyancy force. Thus, the velocity U_c can be obtained by equating the work done by the buoyancy force $g\beta(T_{av} - T_\infty)$ acting along the distance x from the leading edge of the plate with the corresponding change in kinetic energy of the fluid passing from rest at $x = 0$ to U_c at the considered distance x :

$$\frac{1}{2}U_c^2 = g\beta(T_{av} - T_\infty)x \quad (12)$$

where g is the gravitational acceleration and β is the thermal expansion coefficient of the fluid. Hence, assuming that the average temperature is $T_{av} = (T_p + T_\infty)/2$, we have

$$U_c = [g\beta(T_p - T_\infty)x]^{1/2} \quad (13)$$

The same expression for U_c was found by other authors using different approaches—see, e.g., Incropera et al. [9] and Kays et al. [26].

On the other hand, for a given characteristic velocity U_c , the average fluid velocity decreases as the fluid's kinematic viscosity $\nu = \mu/\rho$ increases, since the shear-stress effects, proportional to the fluid's dynamic viscosity μ , become predominant over mechanical inertia, quantified by ρ .

Therefore, the thickness of the thermal boundary layer increases with decreasing U_c and increasing ν .

Notice that the average fluid velocity within the thermal boundary layer tends also to increase with increasing α , due to the greater temperature rise, which implies a larger buoyancy force and a related higher velocity of the upward-moving fluid. Accordingly, since δ_t is inversely proportional to the average fluid velocity within the thermal boundary layer, from this point of view we could conclude that δ_t should decrease with increasing α , which might appear to contradict what we stated earlier. However, an increase in the fluid velocity due to a stronger buoyancy force can only occur if there is a corresponding increase in the fluid temperature and, consequently, if an increase in the thickness of the thermal boundary layer occurs. Therefore, the tendency of δ_t to increase with α must be considered predominant.

Ultimately, for gases and non-metallic liquids, i.e., for $Pr \gtrsim 1$, we can write

$$\delta_t \propto x \left(\text{with decreasing } \frac{d\delta_t}{dx} \right) , \quad \delta_t \propto \alpha , \quad \delta_t \propto 1/U_c , \quad \delta_t \propto \nu \quad (14)$$

Notice that here and throughout the entire paper, the symbol “ $\propto$ ” indicates a qualitative proportionality, while the symbol “ $\sim$ ” is reserved for relations involving a specified functional dependence.

Using Equation (13), the third proportionality relationship of Equation (14) becomes

$$\delta_t \propto \frac{1}{g\beta(T_p - T_\infty)x} \quad (15)$$

where the exponent 1/2 has been omitted, coherently with the meaning of symbol “ $\propto$ ”.

Thus, Equation (14) can be written as

$$\delta_t \propto x \left(\text{with decreasing } \frac{d\delta_t}{dx} \right), \quad \delta_t \propto \alpha, \quad \delta_t \propto \frac{1}{g\beta(T_p - T_\infty)x}, \quad \delta_t \propto \nu \quad (16)$$

Notice that, as far as the dependence on x is concerned, the first and third relationships of Equation (16) may appear contradictory, but this is not true, as the symbol “ $\propto$ ” indicates only proportionality, not an exact functional dependence. Therefore, since the direct proportionality between δ_t and x expressed by the first relation is strongly motivated by the temperature increase in the rising fluid, it is obvious that the inverse proportionality to x included in the third relationship must be weaker than the direct one, so as not to produce any contradiction.

To ensure dimensional consistency and retain as far as possible the proportionalities reported in Equation (16), the ratio δ_t/x must be proportional to one or more dimensionless groups obtained by combining x , α , $g\beta(T_p - T_\infty)x$ and ν , using power laws. So, given that x has dimensions of [m], both α and ν have dimensions of [m²/s] and $g\beta(T_p - T_\infty)x$ has dimensions of [m²/s²], the most natural way to combine these variables into a single dimensionless group is

$$\frac{\delta_t}{x} \propto \frac{\alpha\nu}{[g\beta(T_p - T_\infty)x]x^2} \quad (17)$$

and then

$$\frac{\delta_t}{x} \propto \frac{\alpha\nu}{g\beta(T_p - T_\infty)x^3} \quad (18)$$

where the inverse of the right-hand side of Equation (18) is the Rayleigh number defined in Equation (7). Hence

$$\frac{\delta_t}{x} \propto \frac{1}{Ra} \quad (19)$$

Now, adopting a power law dependence, we can write

$$\frac{\delta_t}{x} \sim \frac{1}{Ra^\gamma} \quad (20)$$

which, according to Equation (7), results in

$$\delta_t \sim \frac{x}{\left[\frac{g\beta(T_p - T_\infty)x^3}{\alpha\nu} \right]^\gamma} \quad (21)$$

Thus, in order to ensure the required direct proportionality between δ_t and x , the positive exponent γ must be lower than 1/3, which also ensures that $d\delta_t/dx$ decreases with increasing x . In this way, all the relationships in Equation (16) are fully satisfied.

At this point, it is useful to emphasize that, in determining Equation (18), the simplest combination of the independent variables was chosen, implicitly assuming that the kinematic viscosity ν and the thermal diffusivity α , which have the same dimensions, give the same exact contribution. This has been proven to be true for high- Pr fluids, for which Equation (20) rigorously holds, yet the attribution of equal importance to ν and α in defining the thickness of the thermal boundary layer is an assumption that, in principle, cannot be taken for granted. In fact, in the previous phenomenological analysis, the two quantities ν and α have been involved in different ways. In this regard, the direct proportionality between δ_t and ν is an implicit dependence resulting from the inverse proportionality between δ_t and the average velocity of the fluid flowing within the thermal boundary layer, which in turn decreases as ν increases. Conversely, the direct proportionality between δ_t and α stems from two opposing contributions: a prevailing explicit direct dependence, since α directly affects the temperature rise in the fluid, leading to a direct proportionality between δ_t and α ; and a weaker implicit inverse dependence, arising from the inverse proportionality between δ_t and the average fluid velocity within the thermal boundary layer, which, as discussed before, increases with increasing α , leading to an inverse proportionality between δ_t and α . Thus, in general, it seems appropriate to include in the relationship expressed by Equation (20) a weak corrective function of the Prandtl number $Pr = \nu/\alpha$ able to account for ν and α having different weights:

$$\frac{\delta_t}{x} \sim \frac{1}{\varphi_1(Pr)Ra^\gamma} \quad (0 < \gamma < 1/3) \quad (22)$$

in which the corrective function $\varphi_1(Pr)$ should be constructed to preserve the direct proportionality between δ_t and both ν and α . It follows

$$\delta_t \sim \frac{x}{\varphi_1(Pr)Ra^\gamma} \quad (0 < \gamma < 1/3) \quad (23)$$

The corresponding heat transfer correlation can be obtained by applying Fourier's law to the fluid film adherent to the plate:

$$h_c(T_p - T_\infty) = -k \left(\frac{\partial T}{\partial z} \right)_0 \quad (24)$$

where h_c is the convection heat transfer coefficient and $(\partial T/\partial z)_0$ is the temperature gradient normal to the plate calculated at the plate surface. Estimating the latter as $(\partial T/\partial z)_0 \sim (T_\infty - T_p)/\delta_t$, we can write

$$h_c \sim \frac{k}{\delta_t} \quad (25)$$

Subsequently, the combination of Equations (23) and (25), taking also into account the definition of the Nusselt number ($Nu = h_c x/k$), leads to the relationship

$$Nu \sim \varphi_1(Pr)Ra^\gamma \quad (0 < \gamma < 1/3) \quad (26)$$

In conclusion, assuming $f_1(Pr) = c_1 \varphi_1(Pr)$, where c_1 is a constant, the following correlation can be written:

$$Nu = f_1(Pr)Ra^\gamma \quad (0 < \gamma < 1/3) \quad (27)$$

which is consistent with Equation (6), in which the exponent γ is equal to $1/4$.

Regarding the structure of the corrective function $\varphi_1(Pr)$, it should be noted that the role played by the Prandtl number in natural convection is significantly different from that

played in forced convection, where it appears raised to the one-third power in the well-known heat transfer correlation found by Pohlhausen [27]. In fact, in forced convection, the Prandtl number serves to relate the thickness δ_t of the thermal boundary layer to the thickness δ of the velocity boundary layer, whose value can be computed regardless of the existence of the temperature field, from which the velocity field is independent—at least as long as the physical properties of the fluid are assumed to be constant. Therefore, given the similarity in the growth of δ_t and δ as the streamwise coordinate x increases, and considering that the descriptive laws for $\delta = \delta(x)$ and $\delta_t = \delta_t(x)$ are power laws, then the function of the Prandtl number expressing δ_t as a function of δ must be a power law as well. Conversely, in natural convection, the Rayleigh number already encompasses all the necessary elements to describe the growth of the thickness δ_t of the thermal boundary layer as a function of the streamwise coordinate x . Thus, it can be reasonably stated that $\varphi_1(Pr)$ plays simply a refining role, intended to account for possible different weights of the kinematic viscosity and the thermal diffusivity in determining the value of δ_t . In this sense, since in natural convection there is a strong and complex interrelation between the temperature and velocity fields, a simple power law could not be adequate to provide a sufficiently accurate representation of the phenomenon's dependence on the Prandtl number, as actually it is not. Of course, as mentioned above, the functional relationship expressed by $\varphi_1(Pr)$ must ensure the direct proportionality of δ_t on ν and α , and thus the inverse proportionality of the Nusselt number on ν and α .

2.2. The Nu-Bo Correlation for Liquid Metals

The results of the previous phenomenological analysis can be extended to liquid metals. Indeed, for low Pr -fluids the fourth relationship in Equation (16) can be neglected, as the kinematic viscosity ν is extremely lower than the thermal diffusivity α . This means that the propagation through the fluid of the friction exerted by the plate is negligible, with the consequence that the average velocity of the fluid within the thermal boundary layer does not depend significantly on ν .

Ultimately, for extremely low- Pr fluids, we can write

$$\delta_t \propto x \left(\text{with decreasing } \frac{\partial \delta_t}{\partial x} \right), \quad \delta_t \propto \alpha, \quad \delta_t \propto \frac{1}{g\beta(T_p - T_\infty)x} \quad (28)$$

Following the same procedure previously used for moderate-to-high- Pr fluids, the individual proportionalities in Equation (28) can be grouped in dimensionless form as

$$\frac{\delta_t}{x} \propto \frac{\alpha^2}{g\beta(T_p - T_\infty)x^3} \quad (29)$$

where the inverse of the right-hand side of Equation (29) is the Boussinesq number, defined in Equation (11).

Hence:

$$\frac{\delta_t}{x} \propto \frac{1}{Bo} \quad (30)$$

Adopting a power law dependence, we can write

$$\frac{\delta_t}{x} \sim \frac{1}{Bo^\varepsilon} \quad (31)$$

where, also in this case, the positive exponent ε must be lower than $1/3$ to ensure both a direct proportionality between δ_t and x and a decreasing $d\delta_t/dx$ with increasing x .

Finally, it should be noted that, in determining Equation (29), the simplest combination of the independent variables was chosen, assuming that the kinematic viscosity ν is completely negligible. In practice, for liquid metals, ν still plays a role, even though it is not dominant over the thermal diffusivity α , especially for relatively high- Pr liquid metals. Thus, Equation (31) can be further refined as

$$\frac{\delta_t}{x} \sim \frac{1}{\varphi_2(Pr)Bo^\varepsilon} \quad (0 < \varepsilon < 1/3) \quad (32)$$

in which the weak corrective function $\varphi_2(Pr)$ should have a structure such to preserve the direct proportionality of δ_t/x from α .

Subsequently, the combination of Equations (32) and (25), taking also into account the definition of the Nusselt number ($Nu = h_c x/k$), leads to

$$Nu \sim \varphi_2(Pr)Bo^\varepsilon \quad (0 < \varepsilon < 1/3) \quad (33)$$

In conclusion, assuming $f_2(Pr) = c_2 \varphi_2(Pr)$, where c_2 is a constant, the following correlation may be written:

$$Nu = f_2(Pr)Bo^\varepsilon \quad (0 < \varepsilon < 1/3) \quad (34)$$

which is consistent with Equation (10), in which the exponent ε is equal to $1/4$. Regarding the structure of $f_2(Pr)$, it should be noted that it is not a power law for the same reasons discussed for gas and non-metallic liquids. Furthermore, as previously said, the functional relationship expressed by $f_2(Pr)$ must ensure the direct proportionality of δ_t on ν and α , and thus the inverse proportionality of the Nusselt number on ν and α .

3. Conclusions

The structure of the heat transfer correlations for laminar natural convection of moderate-to-high- Pr fluids and low- Pr fluids has been examined, using the steady external flow adjacent to an isothermal vertical plate as a reference configuration. A phenomenological boundary layer framework, based exclusively on phenomenological reasoning and deliberately independent of any similarity solution or formal scale analysis, has been employed.

It has been demonstrated that the classical description of natural convection in terms of the Grashof and Prandtl numbers considered separately does not provide a physically effective interpretation of the correlation structure. Instead, the governing dimensionless parameters that naturally emerge are the Rayleigh number for moderate-to-high- Pr fluids and the Boussinesq number for low- Pr fluids, thus obtaining a comprehensive phenomenological description of natural convection flows.

Furthermore, it has been shown that in both cases the exponent of the Rayleigh and Boussinesq numbers must be a positive number lower than $1/3$, and that the Prandtl number enters only through a weak corrective function.

As a supporting result, two appendices have been included providing an overall order-of-magnitude analysis applied to moderate-to-high- Pr fluids and low- Pr fluids to recover the one-fourth power law dependence of the Nusselt number on the Rayleigh and Boussinesq numbers, respectively, at the same time clarifying the emergence of the weak corrective function of the Prandtl number.

This perspective provides a comprehensive physical interpretation of the existing natural convection correlations, also offering guidance for identifying the appropriate dimensionless parameters when developing correlations.

Author Contributions: Conceptualization, M.C., A.Q. and G.D.B.; methodology, M.C., A.Q. and G.D.B.; formal analysis, M.C., A.Q. and G.D.B.; investigation, M.C., A.Q. and G.D.B.; writing—original draft preparation, M.C., A.Q. and G.D.B.; writing—review and editing, M.C., A.Q. and G.D.B.; supervision, M.C. All authors have read and agreed to the published version of the manuscript.

Funding: This research received no external funding.

Institutional Review Board Statement: Not applicable.

Informed Consent Statement: Not applicable.

Data Availability Statement: The original contributions presented in the study are included in the article; further inquiries can be directed to the corresponding author.

Conflicts of Interest: The authors declare no conflicts of interest.

Nomenclature

a	acceleration [m/s^2]
a_b	acceleration due to the only buoyancy force [m/s^2]
Bo	Boussinesq number [–]
c_p	specific heat at constant pressure [$\text{J}/(\text{kg K})$]
F_b	buoyancy force [N]
F_f	friction force [N]
$F(Pr)$	corrective function of the Prandtl number in the $Nu-Gr$ correlation [–]
$f_1(Pr)$	corrective functions of the Prandtl number in the $Nu-Ra$ correlation [–]
$f_2(Pr)$	corrective function of the Prandtl number in the $Nu-Bo$ correlation [–]
g	gravitational acceleration [m/s^2]
Gr	Grashof number [–]
h_c	convection heat transfer coefficient [$\text{W}/(\text{m}^2\text{K})$]
k	thermal conductivity [$\text{W}/(\text{mK})$]
m	mass [kg]
n_1	numerical factor relating u_{av} to U_c [–]
n_2	numerical factor relating u_{max} to U_c [–]
n_3	numerical factor relating δ_u to δ_t [–]
Nu	Nusselt number [–]
Pr	Prandtl number [–]
Q	heat [J]
Ra	Rayleigh number [–]
T_p	plate temperature [K]
T_∞	temperature of the undisturbed fluid [K]
u	vertical velocity [m/s]
U_c	characteristic velocity of natural convection [m/s]
x	vertical coordinate parallel to the plate [m]
W	plate width [m]
z	horizontal coordinate normal to the plate [m]
α	thermal diffusivity [m^2/s]
β	thermal expansion coefficient [$1/\text{K}$]
γ	numerical exponent [–]
ΔT	temperature change [K]
Δt	time interval [s]
δ	thickness of the velocity boundary layer [m]
δ_t	thickness of the thermal boundary layer [m]
δ_u	distance of the velocity maximum from the plate [m]
μ	dynamic viscosity [kg/ms]
ν	kinematic viscosity [m^2/s]
ρ	mass density [kg/m^3]

Subscripts

av	average
0	evaluated at the plate surface ($z = 0$)
max	maximum

Symbols

$\propto$	proportional to (no functional dependence specified)
$\sim$	proportional to (with functional dependence specified)

Appendix A. An Overall Order-of-Magnitude Analysis for Gases and Non-Metallic Liquids

The direct proportionality between the Nusselt and Rayleigh numbers expressed by Equation (27) can be easily derived by an overall order-of-magnitude analysis, whose application, leaving aside any differential formulation of the problem, leads to determine the exact value of the exponent γ .

The analysis is based on the simultaneous application of the heat capacity equation and the Newton's second law of motion to the fluid traveling in laminar flow within the thermal boundary layer from the leading edge of the plate to a generic streamwise position x , as shown in Figure A1, where the profiles of both the temperature and the vertical component of the fluid velocity, and the related thermal and velocity boundary layers, are represented.

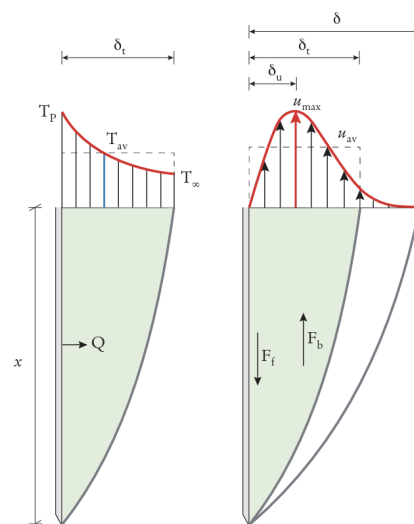


Figure A1. Scheme for the application of the overall order-of-magnitude analysis to gases and non-metallic liquids.

It can be noticed that the velocity profile in the direction normal to the plate originates from zero at the plate surface due to the no-slip condition and asymptotically approaches the zero velocity of the undisturbed fluid far from the plate, necessarily requiring the existence of a maximum at a distance from the plate denoted as δ_u . The distance from the plate where the vertical fluid velocity becomes negligible is defined as the thickness of the velocity boundary layer, denoted as δ . Notice that, due to the tendency of the velocity to reach asymptotically the zero value of the undisturbed fluid, the maximum of the velocity profile is located closer to the plate than toward the edge of the velocity boundary layer, which means that the ratio δ_u/δ is lower than 0.5. Moreover, the maximum of the velocity profile is necessarily located within the thermal boundary layer, where the buoyancy forces act, which means that the ratio δ_u/δ_t does not exceed 1. Finally, it is worth observing that, as ν increases, the average fluid velocity within the thermal boundary layer, u_{av} , decreases and a displacement of the maximum of the velocity profile away from the plate occurs.

The governing equations are:

$$Q = c_p m \Delta T \quad (\text{A1})$$

$$F_b - F_f = ma \quad (\text{A2})$$

where Q is the amount of heat that, having assumed $T_p > T_\infty$, the fluid receives from the considered portion of the plate during the time interval Δt required to reach abscissa x starting from $x = 0$ at the leading edge of the plate, c_p is the fluid's specific heat at constant pressure, m is the mass of fluid involved in the heating process and at the same time subjected to the buoyancy force F_b as well as to the friction force F_f exerted by the considered portion of the plate, ΔT is the temperature change in the mass m of fluid over the time interval Δt , and a is the acceleration consequent to the simultaneous application of F_b and F_f .

Regarding Equation (A1), the heat exchanged between the plate and the fluid can be evaluated by the Fourier's law applied to the motionless fluid film adherent to the plate, which, in the hypothesis that the thermal conductivity k is constant, results in

$$Q = -kWx \left(\frac{\partial T}{\partial z} \right)_{0,av} \Delta t \quad (\text{A3})$$

where W is the plate width and $(\partial T / \partial z)_{0,av}$ is the average temperature gradient in the direction z normal to the plate, evaluated at the plate surface and averaged between $x = 0$ and abscissa x , which can be approximated as

$$\left(\frac{\partial T}{\partial z} \right)_{0,av} \sim \frac{T_\infty - T_p}{\delta_t} \quad (\text{A4})$$

The replacement of Equation (A4) into Equation (A3) results in

$$Q \sim kxW \frac{T_p - T_\infty}{\delta_t} \Delta t \quad (\text{A5})$$

The mass of fluid involved in the heating process can then be evaluated by multiplying the mass flow rate at abscissa x by the duration of the time interval Δt , assuming that the flow is exclusively vertical. Under the assumption of constant mass density ρ , this yields

$$m \sim \rho u_{av} \delta_t W \Delta t \quad (\text{A6})$$

where u_{av} is the average fluid velocity within the thermal boundary layer at abscissa x , expressed as a function of U_c through a scaling factor $n_1 < 1$:

$$u_{av} = n_1 U_c \quad (\text{A7})$$

Notice that, according to what has been previously discussed, n_1 decreases as ν increases and α decreases. Moreover, since n_1 is a pure number and considering that ν and α have the same dimensions $[\text{m}^2/\text{s}]$, n_1 can only decrease with increasing the ratio ν/α , i.e., with increasing the Prandtl number. It then follows

$$m \sim \rho n_1 U_c \delta_t W \Delta t \quad (\text{A8})$$

Finally, the temperature change experienced by the fluid flowing from the leading edge of the plate to abscissa x can be expressed, as

$$\Delta T = T_{av} - T_\infty \quad (\text{A9})$$

where T_{av} denotes the average fluid temperature at abscissa x , which can be evaluated as $T_{av} \sim (T_p + T_\infty)/2$, thus implying

$$\Delta T \sim T_p - T_\infty \quad (\text{A10})$$

The substitution of Equations (A5), (A8) and (A10) into Equation (A1) yields

$$kxW \frac{T_p - T_\infty}{\delta_t} \Delta t \sim c_p \rho n_1 U_c \delta_t W \Delta t (T_p - T_\infty) \quad (\text{A11})$$

which, taking into account the definition of the thermal diffusivity, $\alpha = k/(c_p \rho)$, leads to

$$\delta_t^2 \sim \frac{\alpha x}{n_1 U_c} \quad (\text{A12})$$

As far as Equation (A2) is concerned, the buoyancy force can be determined as

$$F_b = m a_{b,av} \quad (\text{A13})$$

where $a_{b,av} \sim (U_c - 0)/\Delta t$ represents the average acceleration that the fluid would achieve solely due to the buoyancy force when passing from rest at $x = 0$ to U_c at the considered abscissa x . Accordingly, on account of Equation (A8), we have

$$F_b \sim \rho n_1 U_c^2 \delta_t W \quad (\text{A14})$$

On the other hand, in the hypothesis that the dynamic viscosity μ is constant, the friction force can be quantified as

$$F_f = \mu W x \left(\frac{\partial u}{\partial z} \right)_{0,av} \quad (\text{A15})$$

where $(\partial u/\partial z)_{0,av}$ is the average velocity gradient in the direction z normal to the plate, evaluated at the plate surface and averaged between $x = 0$ and abscissa x , which can be approximated as

$$\left(\frac{\partial u}{\partial z} \right)_{0,av} \sim \frac{u_{max} - 0}{\delta_u} \quad (\text{A16})$$

where δ_u is the distance from the plate at which the peak of the velocity profile occurs.

The maximum velocity u_{max} can be assumed to be a fraction of the characteristic velocity U_c , scaled by a factor n_2 , which decreases with increasing the Prandtl number for the same reasons previously discussed for n_1 . In fact, n_1 and n_2 are closely related, since a decrease in the average velocity obviously occurs with a reduction in the maximum velocity. Additionally, δ_u can also be assumed to be a fraction of δ_t scaled by a factor n_3 , which increases with the Prandtl number. As a matter of fact, a higher Prandtl number corresponds to a greater momentum diffusion, which implies an increase in the thickness of the velocity boundary layer and a displacement of the maximum of the velocity profile away from the plate. It follows

$$\left(\frac{\partial u}{\partial z} \right)_{0,av} \sim \frac{n_2 U_c}{n_3 \delta_t} \quad (\text{A17})$$

and then

$$F_f \sim \mu W x \frac{n_2 U_c}{n_3 \delta_t} \quad (\text{A18})$$

With regard to the average acceleration experienced by the mass m during the time interval Δt as it moves from $x = 0$ to a distance x from the leading edge of the plate, we can use the expression

$$a = \frac{\Delta u}{\Delta t} \sim \frac{u_{av} - 0}{\Delta t} \quad (\text{A19})$$

that, on account of Equation (A7), becomes

$$a \sim \frac{n_1 U_c}{\Delta t} \quad (\text{A20})$$

By replacing Equations (A8), (A14), (A18) and (A20) into Equation (A2), we obtain

$$\rho n_1 U_c^2 \delta_t W - \mu W x \frac{n_2 U_c}{n_3 \delta_t} \sim \rho n_1^2 U_c^2 \delta_t W \quad (\text{A21})$$

which implies

$$\rho U_c^2 (n_1^2 - n_1) \delta_t W \sim -\mu W x \frac{n_2 U_c}{n_3 \delta_t} \quad (\text{A22})$$

and then

$$\delta_t^2 \sim \frac{\nu x}{\frac{n_3}{n_2} U_c (n_1 - n_1^2)} \quad (\text{A23})$$

In conclusion, the multiplication of Equations (A12) and (A23), combined with Equation (13), gives

$$\delta_t^4 \sim \frac{\alpha \nu x^2}{\frac{n_3}{n_2} n_1^2 (1 - n_1) g \beta (T_p - T_\infty) x} \quad (\text{A24})$$

Multiplying and dividing the right-hand side of Equation (A24) by x^2 and setting

$$\eta = \frac{n_3}{n_2} n_1^2 (1 - n_1) \quad (\text{A25})$$

we have

$$\delta_t^4 \sim \frac{\alpha \nu x^4}{\eta g \beta (T_p - T_\infty) x^3} \quad (\text{A26})$$

that, on account of Equation (7), can be written as

$$\left(\frac{\delta_t}{x} \right)^4 \sim \frac{1}{\eta Ra} \quad (\text{A27})$$

It follows

$$\frac{\delta_t}{x} \sim \frac{1}{\eta^{1/4} Ra^{1/4}} \quad (\text{A28})$$

At this point, the Nu - Ra relationship can be derived by combining Equations (25) and (A28) together with the definition of the Nusselt number, thus obtaining

$$Nu \sim g_1(Pr) \cdot Ra^{1/4} \quad (\text{A29})$$

where

$$g_1(Pr) = \eta^{1/4} = \left[\frac{n_3}{n_2} n_1^2 (1 - n_1) \right]^{1/4} \quad (\text{A30})$$

which is consistent with Equation (26), as $g_1(Pr)$ represents an estimate of $\varphi_1(Pr)$.

Regarding Equation (A30), the following observations can be made:

- (a) The scaling factor $n_1 = u_{av}/U_c$ is a decreasing function of the Prandtl number. In fact, for an assigned value of the thermal diffusivity α , the average fluid velocity within the thermal boundary layer decreases as the kinematic viscosity ν increases, and then as the Prandtl number Pr increases. This implies that $(1 - n_1)$ is an increasing function of Pr , while n_1^2 is a strongly decreasing function of Pr .
- (b) The scaling factor $n_2 = u_{max}/U_c$ is a decreasing function of the Prandtl number. In fact, for an assigned value of the thermal diffusivity α , the maximum fluid velocity within the thermal boundary layer decreases as the kinematic viscosity ν increases, and then as the Prandtl number Pr increases.
- (c) The scaling factor $n_3 = \delta_u/\delta_t$ is an increasing function of the Prandtl number. In fact, for an assigned value of the thermal diffusivity α , an increase in ν , and then in Pr , reflects a higher aptitude of the fluid to propagate the friction exerted by the plate, which implies a displacement of the maximum of the velocity profile away from the plate.

It then follows that, as the Prandtl number increases, the strong decrease in n_1^2 is somehow counterbalanced by the increase in $(1 - n_1)$, $1/n_2$ and n_3 , making it reasonable to expect that $g_1(Pr) = \eta^{1/4}$ is a weak function of the Prandtl number, which is in line with the literature results.

Appendix B. An Overall Order-of-Magnitude Analysis for Liquid Metals

The direct proportionality between the Nusselt and Boussinesq numbers can be obtained by the same overall order-of-magnitude analysis previously used to derive Equation (A29), by simply neglecting the viscous effects induced by the friction force exerted by the portion of the plate between the leading edge and abscissa x . This means that the overall analysis is based only on the heat capacity equation applied to the fluid traveling in laminar flow within the thermal boundary layer from the leading edge of the plate to a generic streamwise position x . In fact, once the effects of the friction force are neglected, Newton's second law of motion does not provide further information, since it has already been used to derive the expression of U_c reported in Equation (13). A schematic representation of the profiles of both the temperature and the vertical component of the fluid velocity, and the related thermal and velocity boundary layer, is depicted in Figure A2.

In practice, the previous equations (A5), (A6), and (A10) can be replaced into Equation (A1) to obtain an equation that replicates Equation (A12):

$$\delta_t^2 \sim \frac{\alpha x}{n_1 U_c} \quad (\text{A31})$$

Squaring both sides of Equation (A31), we obtain

$$\delta_t^4 \sim \frac{\alpha^2 x^2}{n_1^2 U_c^2} \quad (\text{A32})$$

whose combination with Equation (13) gives

$$\delta_t^4 \sim \frac{\alpha^2 x^2}{n_1^2 g \beta (T_p - T_\infty) x} \quad (\text{A33})$$

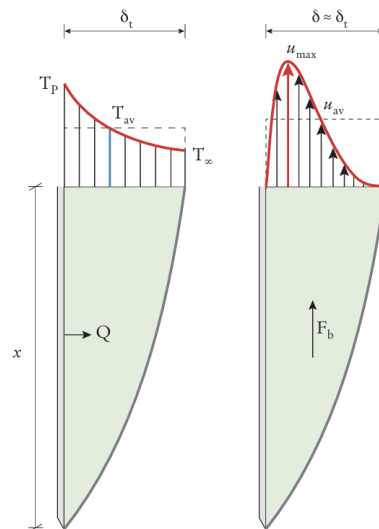


Figure A2. Scheme for the application of the overall order-of-magnitude analysis to liquid metals.

Multiplying and dividing the right-hand side of Equation (A33) by x^2 , we have

$$\delta_t^4 \sim \frac{\alpha^2 x^4}{n_1^2 g \beta (T_p - T_\infty) x^3} \quad (\text{A34})$$

that, on account of Equation (11), can be written as

$$\left(\frac{\delta_t}{x}\right)^4 \sim \frac{1}{n_1^2 Bo} \quad (\text{A35})$$

It follows

$$\frac{\delta_t}{x} \sim \frac{1}{n_1^{1/2} Bo^{1/4}} \quad (\text{A36})$$

At this stage, the Nu - Bo relationship can be derived by combining Equations (25) and (A36), together with the definition of the Nusselt number, thus obtaining

$$Nu \sim g_2(Pr) \cdot Bo^{1/4} \quad (\text{A37})$$

where

$$g_2(Pr) = n_1^{1/2}. \quad (\text{A38})$$

which is consistent with Equation (33) as $g_2(Pr)$ represents an estimate of $\varphi_2(Pr)$. Notice that the scaling factor $n_1 = u_{av}/U_c$ is a decreasing function of the Prandtl number. In fact, for an assigned value of the thermal diffusivity α , the average fluid velocity within the thermal boundary layer decreases as the kinematic viscosity ν increases, and then as the Prandtl number Pr increases. On the other hand, since for all low- Pr fluids, the kinematic viscosity ν , and thus the friction force, is almost negligible, it follows that the velocity profile directly arises from the buoyancy force profile, which basically reflects the temperature profile. As a consequence, the average fluid velocity within the thermal boundary layer is not dissimilar from the characteristic velocity U_c calculated considering the fluid at the average temperature T_{av} , as reported in Equation (12). Consequently, n_1 , and so $g_2(Pr)$, is a weak function of Pr , moderately smaller than unity, whatever the value of the Prandtl number in the typical range of liquid metals.

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