A new model for Double Diffusion + Turbulence

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Available models of Double Diffusion (DD) processes (Salt Fingers, SF and Diffusive Convection, DC) are primarily based on laboratory experiments that do not include turbulence which is however always present in the ocean. A reliable DD model for use in OGCMs (ocean global circulation models) is therefore still lacking and a true assessment of the role of oceanic DD is yet to be made. Here, we derive and validate a new model for DD + Turbulence using a second-order closure model which differs significantly from previous ones in that the ratios of the correlation time scales to the dissipation time scales, constant in previous models, now depend on Ri and R_p, the key new feature needed to reproduce both laboratory (no shear) and oceanic (with shear) data. The model can therefore be used in OGCMs. The full mixing model includes mixed layer, internal gravity waves, DD and tides.


1. Introduction

Maps by D. E. Kelley (personal communication, 2007) of the density ratio R_p = (a_1 \partial S/\partial z)(a_1 \partial T/\partial z)^{-1} (T and S are mean temperature and salinity fields and \partial_T S are the thermal expansion and haline contraction coefficients) show the ocean's regions that are prone to DD instabilities [Ruddick and Gargett, 2003]. From those maps one observes that the likelihood of SF is higher in the Atlantic than in most of the Pacific and that DC may play a significant role in the Arctic and in the Southern Ocean, a point discussed by Kelley et al. [2003] who concluded that DC "could be of major importance to the properties of the global ocean". DC is more likely in high-latitude precipitation zones [Schmitt, 1994] and Muench et al. [1990] also found it over much of the Weddell Sea. Overall, in the circumpolar current, both SF and DC may be important. These high latitude regions are of dynamical interest since numerical simulations [Hasumi and Suginohara, 1999; Webb and Suginohara, 2001] show that a sizeable upwelling of the thermohaline circulation occurs in those regions, in contrast to the traditional view of a uniform upwelling throughout the whole ocean. Since thus far most DD models have been based on laboratory data [see Kunze, 2003; Schmitt, 1994, 2003; Kelley et al., 2003], they do not represent the true oceanic environment where there is always a turbulent background [the latter can be represented by a Richardson number Ri = N^2/\Sigma^2, where N is the Brunt-Vaisala frequency (N^2 = \partial b/\partial z, b = -g/\rho_0 is the buoyancy); \Sigma is the mean shear (\Sigma^2 = (\partial U/\partial z)^2 + (\partial V/\partial z)^2 where U, V are the horizontal mean velocity components)]. Models for DD + Turbulence for arbitrary Ri and R_p have been proposed [Walsh and Ruddick, 2000; Inoue et al., 2007] but the presence of unknown parameters precludes their use in OGCMs. Smyth and Kimura [2007] employed a linear stability analysis to study DD + shear but the heat mixing efficiency \Gamma_h vs. R_p was opposite to that of the data (see section 6).

In this paper we present a new model for DD + Turbulence valid for arbitrary Ri, R_p and test it against laboratory data (R_p very large Ri) and oceanic data (large range of values of Ri and R_p). The data are \gamma (heat to salt flux ratio) for SF, R_F (salt to heat flux ratio) for DC and \Gamma_h (heat mixing efficiency) for SF:

\gamma(R_i, R_p), R_F(R_i, R_p), \Gamma_h(R_i, R_p)

Specifically, \gamma(R_i, R_p) in the Ri \gg 1 limit is taken from St. Laurent and Schmitt [1999], Kunze [2003] and Schmitt [2003] \Gamma_h(R_i, R_p) is taken for Ri \gg 1 from Kelley [1990] and R_F(R_i, R_p) is taken from St. Laurent and Schmitt [1999]. Since the model reproduces the functions in equation (1) reasonably well, it can be used in OGCMs which need heat and salt diffusivities K_{h,s} that depend on the variables (1).

2. Heat and Salt Diffusivities

Though we are primarily interested in heat and salt fluxes, their dynamic equations depend on other second-order moments and one must therefore consider \partial_u \partial S, \partial_u \partial T, \partial S, \partial^2 S, \partial T, \partial^2 T, representing momentum, heat and salt fluxes, temperature and salinity variances and temperature-salinity correlation whose dynamic equations were presented in equations (5)–(11) of Canuto et al. [2002] (hereinafter referred to as C2). In the local and stationary case, such equations can be solved analytically with the following results:

\begin{align}
\partial_u \partial S &= -K_\alpha \partial U/\partial z, \quad \partial_u \partial T &= -K_\alpha \partial T/\partial z, \quad \partial S &= -K_\alpha \partial S/\partial z, \quad \partial b &= -K_\alpha N^2 \\
K_\alpha &= \Gamma_\alpha \frac{\varepsilon}{N^2}, \quad \Gamma_\alpha = \frac{1}{2} \left( \gamma N^2 \right)^2 S_u, \quad \Gamma_p = \Gamma_h (1 - \gamma^{-1}) (1 - R_p)^{-1}
\end{align}

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where \( i = 1, 2 \) (\( \alpha \) stands for momentum, heat, salt and density); \( \tau = 2K/\varepsilon \) is the dynamical time scale, \( K \) is turbulent kinetic energy and \( \varepsilon \) is its rate of dissipation. The form of the dimensionless structure functions \( S_h(Ri, R_p) \) is given by equations (13a)–(15) of C2. Here, we present an alternative form for \( S_{h,s} \) in terms of the ratio \( \bar{w}^2/K \):

\[
S_h = A_h \left( \frac{\bar{w}^2}{K} \right),
\]

\[
S_s = A_s \left( \frac{\bar{w}^2}{K} \right),
\]

\[
\frac{\bar{w}^2}{K} = 2 - \frac{4}{15} \frac{3 \Gamma_n}{\gamma(Ri + R_p)}
\]

\[
A_h = \pi_4 \left[ 1 + px + \pi_2 \pi_4 x (1 - \gamma^{-1}) \right]^{-1},
\]

\[
A_s = \pi_1 \left[ 1 + qx + x \pi_1 \pi_3 R_p (\gamma - 1) \right]^{-1}
\]

where \( \gamma \) is the heat to salt flux ratio given by:

\[
\gamma = \frac{\alpha_T \bar{w}^2}{\alpha_S \bar{w}^2} = \frac{K_h}{K_s R_p} = \frac{1}{\pi_1} \frac{1 + qx}{1 + px}
\]

where

\[
q = \pi_1 \pi_5 (1 + R_p) - \pi_1 \pi_3 R_p, \quad p = \pi_4 \pi_5 - \pi_2 \pi_4 (1 + R_p), \quad x = (\tau N)^2 (1 - R_p)^{-1}
\]

3. Dynamical Time Scale \( \tau \) and the Variable \( x \)

Here we discuss how to determine the dynamical time scale \( \tau \), equation (2b), or \( x = (\tau N)^2 (1 - R_p)^{-1} \), equation (2e).

Consider the production = dissipation relation:

\[
P_m + P_b = \varepsilon \Rightarrow K_m \Sigma^2 - K_p N^2 = \varepsilon \Rightarrow \Gamma_{m}^{-1} \Gamma_m - \Gamma_p = 1 \tag{3a}
\]

where \( P_m \) is the shear production defined as usual as:

\[
P_m = -\left( \bar{w}^2 U_x + \bar{w}^2 V_x \right) = K_m \Sigma^2 \tag{3b}
\]

If in the third relation in (3a), which follows from the second using the first of (2b), we substitute (2b, e) and use the dimensionless structure function \( S_m(x, Ri, R_p) \) given by equation (13a) of C2, we obtain an algebraic relation for \( x \) the solution of which is:

\[
x = x(Ri, R_p, \gamma, \pi s x) \Rightarrow x(Ri, R_p) \tag{3c}
\]

The last step comes from the relations for the \( \pi \)’s (\( Ri, R_p \)) derived in the next section.

4. Key New Ingredients: The Relaxation-Dissipation Time Scales

The dynamic equations for the five second-order moments discussed before contain relaxation-dissipation time scales, which in dimensionless form, are called:

\[
\pi_1 = \tau_{p_\theta} / \tau, \quad \pi_2 = \tau_{\theta_4}/\tau, \quad \pi_3 = \tau_s/\tau, \quad \pi_4 = \tau_{p_\theta}/\tau, \quad \pi_5 = \tau_s/\tau
\]

which were traditionally [e.g., Mellor and Yamada, 1982] assumed constant. An improved set of constants were derived by C2. The C2 values, denoted by a superscript zero, are: \( \pi_1^0 = \pi_4^0 = (27K_0/2)^{1/2}(1 + \sigma_1^{-1})^{-1}, \quad \pi_2^0 = \pi_3^0 = \sigma_1, \quad \pi_5^0 = 1/3 \). Here, \( K_0 \) is the Kolmogorov constant and \( \sigma_1 = 0.72 \) is the (neutral) turbulent Prandtl number. Such relations do not provide a good fit to the variables (1), see Figures 1, 2 and 3 We therefore had to construct a new model in which the \( \pi \)’s are functions of the two key variables:

\[
\pi_k(Ri, R_p) \tag{4b}
\]

The starting point was a recent study [Canuto et al., 2007] (hereinafter referred to as C7) [see also Zilitinkevich et al., 2007] that extended the traditional second-order closure models [e.g., Cheng et al., 2002; C2] to accommodate a new set of DNS (direct numerical simulations), LES (large eddy simulations), lab and field data etc. that showed that mixing persists at almost any Ri dispensing the traditional notion that there exists a critical Ri(cr) above which turbulent mixing essentially vanishes. Specifically, it was shown that the most crucial time scale is \( \tau_{p_\theta} \) that enters the
heat flux equation. The underlying idea is that stable stratification reduces the heat flux more than the momentum flux [Gerz et al., 1989], thus pointing toward $t_p$ as a key player. In fact, a dependence on $R_i$ means the interplay between the temperature and velocity fields and the combination most likely to be affected is the heat flux $w_q$ (thus $t_p$) that entails both those fields. A second reason is the work by Weinstock [1978] who showed that stable stratification reduces the time scale by a factor $1 + (t N)^2$ which C7 showed to be equivalent to:

$$t_p = \frac{t}{C_24} + \frac{R_i}{C_0}$$

To include the effect of both $R_i$ and $R_r$, we now suggest to further generalize (4c) to:

$$t_p = \frac{t}{C_24} + \frac{R_i + aR_r}{C_0}$$

(4d)

For $R_r < 0$, there is no DD tendency and we use (4c). Why the $R_r$ dependence and what is $a$? As for $R_r$, the rationale is as follows: in the presence of DD, for $R_r > 0$, when the salt and temperature gradients act against each other, DD provides a source of mixing in addition to that of shear thus lessening the damping effect of $R_i$, a fact reflected by the counter factor $(1 + aR_r)$. In the limit $R_r \gg R_i$, $\pi_1 \sim \pi_4^0$, the effect of $R_i$ on the time scale disappears. As for $\pi_1$, its form is derived from (4d) by the symmetry requirement:

$$\tau_{ps} \sim \tau_{\rho \theta} \text{ as } R_\rho \to R_\rho^{-1}$$

(4e)

and thus we have:

$$\tau_{ps} / \tau \sim \left[1 + Ri \left(1 + aR_\rho^{-1}\right)^{-1}\right]^{-1}$$

(4f)

Note that in the limit of no temperature gradient, $R_\rho = \infty$, $\pi_1 \sim \pi_1^0(1 + Ri)^{-1}$ in analogy with (4c). As for $\pi_2$, we suggest a generalization that does not depend on $R_i$ since this time scale represents a correlation between temperature and salinity fields only:

$$\tau_{\theta \theta} / \tau \sim \left(R_\rho + R_\rho^{-1}\right)^{-1}$$

(4g)

because of the following reasons. First, it is the simplest form that satisfies the symmetry requirement:

$$\tau_{\theta \theta} \sim \tau_{\theta \theta} \text{ as } R_\rho \to R_\rho^{-1}$$

(4h)
Second, as $R_p \sim 0$, $\pi_2 \sim R_p$ and $\pi_0 \sim 1$ given by (2f) yields the results represented by high Ri. For the case with tendency to DC ($1 < R_p < \infty$), $R_p^{-1}$ plays the same role that $R_i$ does in the case with tendency to SF ($0 < R_p < 1$), when either SF disappears ($R_p = 0$) or DC disappears ($R_p = \infty$), the salt-heat correlation time scale goes to zero, as these correlations are associated with the double diffusive processes. By heat-salt symmetry, $R_p$ and $R_i^{-1}$ must appear symmetrically in the salt-heat correlation. Their average, representing double diffusive tendency of either kind, plays a role in the salt-heat correlation somewhat similar to that of stratification in the pressure correlations.

In summary, $\pi_{1,2,4}$ were modified as just described while $\pi_{3,5} = \pi_{3,5}^0$. Finally, we discuss the coefficient $a$. Consider DC in the limit $R_i \gg 1$ where the data are given in Figure 2. In that limit, the asymptotic values of $\pi_{1,4}$ are:

$$
\pi_4 = a \pi_1^0 R_p R_i^{-1}, \quad \pi_1 = a \pi_1^0 R_i^{-1}
$$

(4i)

It follows that, using (2e), the salt to heat flux ratio $R_F$ becomes:

$$
R_F = \frac{\alpha_b w_S}{\alpha_T w_{\theta}} \Rightarrow 1 \frac{\pi_1^0}{\pi_4} = \frac{\pi_1}{a} \Rightarrow 1
$$

(4j)

Next, using Linden’s [1974] result that the asymptotic value of $R_F$ is the square root of the ratio of the salt diffusivity $\kappa_s$ to the thermal diffusivity $\kappa_T$, we obtain:

$$
a = (\kappa_T/\kappa_s)^{1/2} \approx 10
$$

(4k)

which is the value used in our work. In summary, we now have the new relations:

$$
\begin{align*}
\pi_1 &= \pi_1^0 \left(1 + \frac{R_i}{1 + aR_p}\right)^{-1}, \\
\pi_4 &= \pi_4^0 \left(1 + \frac{R_i}{1 + aR_p}\right)^{-1}, \\
\pi_2 &= \pi_2^0 \left[\frac{1}{2} \left(R_p + R_i^{-1}\right)\right]^{-1}, \quad \pi_{3,5} = \pi_{3,5}^0
\end{align*}
$$

(4l)

which yield the $\pi$’s in terms of $R_i$ and $R_p$.

5. Full Model

The mixing model is thus complete. First, with the $\pi_i(R_i, R_p)$ of (4l) and (3c), one computes $\gamma_i(R_i, R_p)$, equation (2e). Next, with $\gamma_i(R_i, R_p), \pi_i(R_i, R_p)$ and $x(R_i, R_p)$, from (2b–e) one derives the diffusivities of momentum, heat, salt and density as functions of $R_i$ and $R_p$. Here, we do not need to specify $\varepsilon \varepsilon$ since we only deal with flux ratios and mixing efficiencies. However, in an OGCM, one must employ the relation:

$$
\varepsilon = \varepsilon(ML) + \varepsilon(igw) + \varepsilon(tides)
$$

(5)

$\varepsilon(ML)$ and $\varepsilon(igw)$ were discussed in C2, while $\varepsilon(tides)$ is given by Jayne and St. Laurent [2001].

6. Model Results Vs. Data

In Figure 1 we show the data of heat to salt flux ratio $\gamma(R_i, R_p)$ vs. $R_p$ for the SF case (see St. Laurent and Schmitt [1999] for references on data) on which we have superimposed the model results for $R_i = 5, 10^4$ for two different models of the $\pi$’s. In the $\pi_i^0$ case, the model’s predictions (short-dashed lines), bunch up in the upper left corner outside the bulk of the data. Furthermore, the model becomes unrealizable for $R_p \leq 1/2$. With the $\pi$’s from (4l), the model results (solid and long-dashed lines) reproduce the data for large values of $R_i$ (lab data). Model results for $R_p > 10^4$ are indistinguishable from those at $R_i = 10^4$.

In Figure 2 for DC, we show the salt to heat flux ratio $R_F$ vs. $R_p$ also for the two models of the $\pi$’s. As in Figure 1, the $\pi_i^0$ case yields results (short-dashed lines) that bunch up in the upper left corner outside the bulk of the data and with the incorrect $R_p$ dependence. Moreover, the model becomes unrealizable for $R_p \geq 2$. Dots and crosses represent the data with no shear by Kelley [1990] who also suggested an empirical fitting formula (solid line). With the $\pi$’s from (4l), the results reproduce the data much better showing the importance of the $R_i$ and $R_p$ dependence. As in Figure 1, the model results at high $R_i$ reproduce the no-shear data better.

In Figure 3 we show the heat mixing efficiency $\Gamma_h(R_i, R_p)$ using $\pi_i^0$. The model results (dashed and full lines) are superimposed on the data from St. Laurent and Schmitt [1999]. The model results are acceptable for the strong turbulence case ($R_i < 1$), panels a)–c), but not as
good for the remaining cases. In Figure 4, we show the heat mixing efficiency $G_h(R_i, R_{tr})$ using equation (4l). The model results (dashed and full lines) are superimposed on the same data as in Figure 3. In Figure 4f, the model results for $R_i > 10^{4}$ are indistinguishable from those for $R_i = 10^{4}$. Use of (4l) considerably improves the model’s predictions. As in Figures 1 and 2, as $R_i \rightarrow \infty$, the model results become independent of $R_i$ when DD is active.

7. Conclusions

[13] A model has now been constructed which reproduces reasonably well both laboratory DD data ($R_i \gg 1$) and ocean data for a range of $R_i$. The value $\Gamma_h = 0.2$ is now seen to be limited to regions of strong turbulence and no DD, first panel in Figure 3. However, when $R_i \geq 1$ and turbulence still exists but is weak, DD is active and produces $\Gamma_h$’s up to three times as large as 0.2, which casts doubts on the computations that employ 0.2 to discuss the “closure of the thermohaline circulation”. A map of the density mixing efficiency $\Gamma_\rho$ showing regions of positive and negative buoyancy, is needed to visualize where DD processes are relevant: if $\Gamma_\rho < 0$, buoyancy acts like a source of mixing whereas in regions where $\Gamma_\rho > 0$ buoyancy acts like a sink, as in ordinary stably stratified flows without DD. As a concrete example, consider the advection-diffusion model where $\mathbf{\mathcal{W}}$ is the diapycnal advection velocity:

$$\mathbf{\mathcal{W}} N^2 = K_\rho \frac{\partial N^2}{\partial z} + N^2 \frac{\partial K_\rho}{\partial z} \quad (6a)$$

Though stratification increases from the bottom up ($\partial N^2/\partial z > 0$), this does not mean that the first term implies upwelling $\mathbf{\mathcal{W}} > 0$ since regions of strong DD entail $K_\rho < 0$ and thus the first term can lead to downwelling $\mathbf{\mathcal{W}} < 0$. A further important variable is the buoyancy flux integrated over a control volume between two specified density surfaces (TW = terawatts):

$$P_b(TW) = - \int \rho K_\rho N^2 dV = - \int \rho \Gamma_\rho dV \quad (6b)$$

When $P_b > 0$, corresponding to a buoyancy gain, the result can be compared with other sources of ocean stirring such as
as wind and tides [Wunsch, 2000]. With the availability of 
the new model, these studies are presently being pursued.

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