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Armand Vic, Xavier Carton, Jonathan Gula. Eady baroclinic instability of a circular vortex. *Symmetry*, 2022, 14 (7), pp.1438. 10.3390/sym14071438 . hal-03926089

HAL Id: hal-03926089

<https://cnrs.hal.science/hal-03926089>

Submitted on 6 Jan 2023

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EADY BAROCLINIC INSTABILITY OF A CIRCULAR VORTEX

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Abstract: The stability of two superposed buoyancy vortices is studied linearly in a two-level SQG model. The basic state is chosen as two top-hat vortices (with uniform buoyancy), coaxial and with same radius. Only the vertical distance between the two levels and the top and bottom buoyancy intensities are varied, the other parameters are fixed. The linear perturbation equations around this basic state form a two-dimensional ODE for which the normal and singular mode solutions are numerically computed. For normal modes, the system is stable if the vortices are sufficiently far from the other to prevent vertical interactions of the buoyancy patches, or if they are close to each other but with very different intensities, again preventing the resonance of Rossby waves around their contours. The vortex is unstable if the intensities are similar and if the vortices are close to each other vertically. The growth rates of the normal modes increase with the angular wave-number, also corresponding to shorter vertical distances. The growth rates of the singular modes do not depend much on the bottom buoyancy at short time, but, as expected, they converge towards the growth rates of the normal modes. This study remaining linear does not predict the final evolution of such unstable vortices. This nonlinear evolution will be studied in a sequel of this work.

Keywords: SQG; Rankine vortices; normal modes; singular modes; linearisation around basic state.

Citation: Vic, A.; Carton, X.; Gula, J. Eady vortex problem. *Symmetry* **2021**, *1*, 0. <https://doi.org/>

Received: XXX

Accepted: XXX

Published: XXX

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1. Introduction

Vortices are energetic features in many turbulent flows. Paramount among them are geophysical flows, where vortices play an essential role in the planetary transport of energy, heat, moisture for atmospheric vortices, and salinity for oceanic vortices. Vortices are long-lived recirculation motion with a lifetime longer than their turnover period. Therefore it is essential to study the mechanisms underlying their robustness, or their possible destabilization. This has been the subject of many papers in the past ([1–5] and references therein). But many of these papers considered vortex instability in a layered model of the ocean (often corresponding to the ocean above the main thermocline, i.e. above 500 m depth, and below it). This problem is then called the Phillips baroclinic instability of these vortices. The focus was then on fairly large vortices in the ocean (vortices wider than 30 km in radius). Far fewer papers were devoted to the study of vortex stability in a level model of the ocean, where only the density interfaces are concerned. This problem is called the Eady baroclinic instability of these vortices. Such interfaces (the ocean surface, the ocean bottom or the thermocline) play an essential role in ocean dynamics. It was shown recently that a model describing these surfaces only, can represent the dynamics of smaller vortices (with radii 10–30 km) which are abundant in the ocean. Such a model is the Surface Quasi-Geostrophic (SQG) model, employed in this study. This model describes the time evolution of buoyancy anomalies on surfaces, in a rotating stratified flow, with null internal potential vorticity. The 3D internal dynam-

38 ics (vertically, between the horizontal surfaces) are driven by the buoyancy anomalies
 39 on these surfaces.

40 The previous studies of vortex stability in the SQG model concerned the horizontal
 41 shear (barotropic) instability of a single vortex in one- or two-level configurations. The
 42 analytical stability of two superposed vortices in a two-level SQG model has not been
 43 investigated before. Badin and Poulin [6] or Harvey and Ambaum [7] studied the
 44 barotropic instability of a single vortex. Here, we study the baroclinic instability (ver-
 45 tical shear instability of a rotating fluid) of two superimposed vortices. The two-level
 46 SQG model (see [8]) is adapted to deal analytically with this vortex instability problem.
 47 Analytically, the model solves a hyperbolic equation for the transport of buoyancy of the
 48 surfaces and a 3D Laplacian equation on the streamfunction (an elliptical equation to
 49 invert the buoyancy distribution into a flow field) with Neumann boundary condition.
 50 The model equations and their numerical implementation are developed in section 2.
 51 Section 3 presents the basic state composed of two vortices (one at the ocean surface, one
 52 at the ocean thermocline or bottom). These top-hat vortices can have different intensities
 53 and their vertical separation can be varied. In section 4, we linearize the equations
 54 around this steady state and we study how the perturbation grows. This perturbation
 55 can be a normal mode or a singular mode. Section 4 presents and interprets the results.
 56 A conclusion and perspectives follow. Two appendices present details about analytical
 57 and numerical computations.

58

59 2. Surface quasi-geostrophic model and equations

60 The framework is a two-level surface quasi-geostrophic (SQG) model. The two
 61 horizontal surfaces are the actual surface and the bottom of the ocean (or the thermocline
 62 if the temperature - and density - gradient is sufficiently abrupt at this depth [9,10]);
 63 they are vertically separated by a height H (see Figure 1). The buoyancy (or potential
 64 temperature) distributions are contained in the two levels. They are connected by a
 65 condition of null potential vorticity inside the ocean.

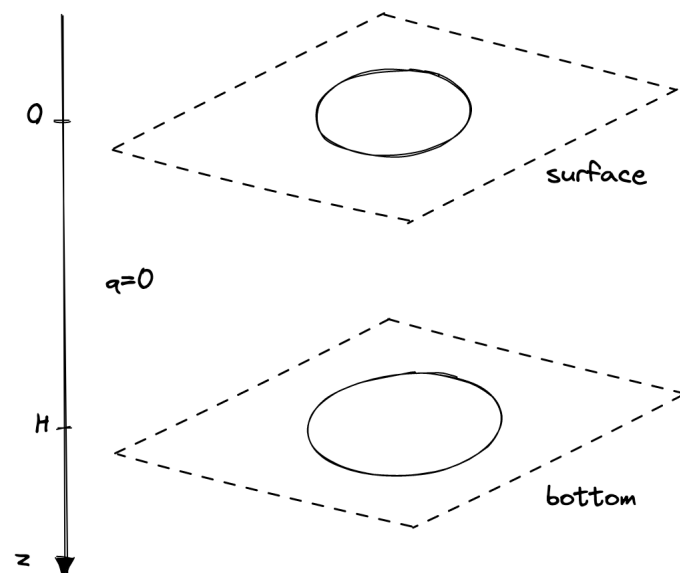


Figure 1. Scheme of the two horizontal layers.

The surface quasi-geostrophic (SQG) model is the restriction of the complete quasi-geostrophic model – introduced by Charney in 1948[11] – to null internal potential vorticity distributions (for more details on the SQG model, see [8,12]). Potential vorticity is thus concentrated as a vertical Dirac distribution, reducing to a planar buoyancy

anomaly, at the two (upper and lower) boundaries.

Assuming constant Brunt-Väisala and Coriolis frequencies, the quasi-geostrophic model is governed by the conservation of potential vorticity q in the fluid volume, in the absence of forcing and of dissipation for the flow:

$$dq/dt = \partial_t q + J(\psi, q) = 0, \quad \text{for } 0 < z < H,$$

associated with following 3D Laplace equation :

$$q = \nabla_h^2 \psi + \frac{f_0^2}{N_0^2} \partial_z^2 \psi, \quad \text{for } 0 < z < H, \quad (1)$$

where J is the horizontal Jacobian operator and ψ is the streamfunction (remember that the horizontal velocity is $u = -\partial_y \psi, v = \partial_x \psi$). The boundary conditions of the 3D model at the top and bottom of the domain, are the horizontal advection of the buoyancies b on these surfaces:

$$[\partial_t + J(\psi, \cdot)]b = 0, \quad z = 0, H \quad (2)$$

66 with $b = f_0 \partial_z \psi$.

67 The surface quasi-geostrophic equations are therefore the restriction of this model to
68 $q = 0$ in the fluid interior. This leads to a model defined only in terms of the surface and
69 bottom buoyancies, related to streamfunction as above.

70 From these equations, one can define a horizontal length scale, called the internal radius
71 of deformation, via $\sigma = \frac{N_0 H}{f_0}$. This scale represents the distance over which the buoyancy
72 and Coriolis accelerations have similar intensities. Now we normalize in $[0, 1]$ the vertical
73 scale (H) to have the following set of equations :

$$\begin{cases} \nabla_h^2 \psi + \frac{1}{\sigma^2} \partial_z^2 \psi = 0 & \text{for } 0 < z < 1 \\ \left. \frac{\partial \psi}{\partial z} \right|_{z=0} = b^s, \quad \frac{D b^s}{D t} = 0 \\ \left. \frac{\partial \psi}{\partial z} \right|_{z=1} = b^b, \quad \frac{D b^b}{D t} = 0 \end{cases} \quad (3)$$

74 where $\frac{D}{D t}$ is the horizontal Lagrangian derivative and the superscripts s and b represent
75 respectively "surface" and "bottom".

76

The first equation of system (3) in horizontal Fourier space (k, l, z) gives

$$\frac{\partial^2 \hat{\psi}}{\partial z^2} = K^2 \sigma^2 \hat{\psi} \quad (4)$$

77 where $K^2 = k^2 + l^2$. This equation with boundary conditions for buoyancies gives in
78 Fourier space

$$\hat{\psi}(k, l, z) = \frac{1}{K\sigma \sinh(K\sigma)} \left(\hat{b}^b \cosh(K\sigma z) - \hat{b}^s \cosh(K\sigma(1-z)) \right) \quad (5)$$

79 The mean flow in our SQG model are two top-hat vortices (i.e. two vortices with
80 constant buoyancy in a disk of radius unity). Here we analyse the linear stability of this
81 mean flow (or basic state). This problem is called the Eady baroclinic instability of this
82 vortex. With this mean flow geometry, the polar coordinates are a natural choice.

83 3. Mean flow calculation

84 Here we calculate the flow field associated with these two top-hat vortices. Firstly
85 we remind the form of cylindrical Fourier transforms.

86 3.1. Preliminaries about Fourier decomposition in cylindrical coordinates

87 In cylindrical coordinates, consider a function $f(r, \phi, z)$ sufficiently regular.

- Then f is 2π -periodic in ϕ so can be decomposed in Fourier modes :

$$f(r, \phi, z) = \sum_{n \in \mathbf{N}} \tilde{f}(r, n, z) e^{in\phi} \quad (6)$$

with

$$\tilde{f}(r, n, z) = \frac{1}{2\pi} \int_0^{2\pi} f(r, \phi, z) e^{-in\phi} d\phi \quad (7)$$

- For every $n \in \mathbf{N}$ and $z \in \mathbf{R}^+$, $\tilde{f}(\cdot, n, z)$ are functions which can be written as inverse Hankel transforms :

$$\tilde{f}(r, n, z) = \int_0^\infty \hat{f}(\rho, n, z) J_n(\rho r) \rho d\rho \quad (8)$$

where J_n are the Bessel functions and with

$$\hat{f}(\rho, n, z) = \int_0^\infty \tilde{f}(r, n, z) J_n(\rho r) r dr \quad (9)$$

In fine, the function can be decomposed as

$$f(r, \phi, z) = \sum_{n \in \mathbf{N}} \int_0^\infty \hat{f}(\rho, n, z) J_n(\rho r) \rho d\rho e^{in\phi} \quad (10)$$

with

$$\hat{f}(\rho, n, z) = \frac{1}{2\pi} \int_0^{2\pi} \int_0^\infty f(r, \phi, z) J_n(\rho r) r dr e^{-in\phi} d\phi \quad (11)$$

88 where \hat{f} is the Fourier transform of f in cylindrical Fourier coordinates.

89 **Remark 1.** For the second bullet point, the Bessel function could be arbitrary but the n -th
90 function is retained because it is a solution of the Laplace equation.

91 3.2. Application to SQG flows

92 We decompose ψ as in (10) to get :

$$\psi(r, z, \phi) = \sum_{n \in \mathbf{N}} \int_0^\infty \hat{\psi}(\rho, n, z) J_n(\rho r) \rho d\rho e^{in\phi} \quad (12)$$

93 with $\hat{\psi}$ the horizontal Fourier transform of ψ , already computed in (5). We deduce :

$$\psi = \sum_{n \in \mathbf{N}} \int_0^\infty \frac{\rho J_n(\rho r)}{\rho \sigma \sinh(\rho \sigma)} \left(\hat{b}^b(\rho, n, t) \cosh(\rho \sigma z) - \hat{b}^s(\rho, n, t) \cosh(\rho \sigma (1 - z)) \right) d\rho e^{in\phi}. \quad (13)$$

94 So at the two boundaries (surface and bottom) we have :

$$\psi^s(r, \phi, z = 0, t) = \sum_{n \in \mathbf{N}} \int_0^\infty \frac{J_n(\rho r)}{\sigma \sinh(\rho \sigma)} \left(\hat{b}^b - \hat{b}^s \cosh(\rho \sigma) \right) d\rho e^{in\phi} \quad (14a)$$

$$\psi^b(r, \phi, z = 1, t) = \sum_{n \in \mathbf{N}} \int_0^\infty \frac{J_n(\rho r)}{\sigma \sinh(\rho \sigma)} \left(\hat{b}^b \cosh(\rho \sigma) - \hat{b}^s \right) d\rho e^{in\phi} \quad (14b)$$

95 **Remark 2.** From now, we will denote by capital letters the basic state variables and by lowercase
96 letters the perturbed variables. For example, the total streamfunction at the surface will be
97 $\Psi^s + \psi^s$.

98 3.3. Basic state : two top-hat vortices

99 We take as basic states two top-hat vortices (i.e two circular plateaus of constant
100 buoyancy) with the same dimensionless radius $R = 1$ but not the same intensity.

$$B^s = B_0^s(1 - H(r - 1)) \quad (15a)$$

$$B^b = B_0^b(1 - H(r - 1)) \quad (15b)$$

101 where $H(x) = \begin{cases} 0 & \text{if } x < 0 \\ 1 & \text{if } x > 0 \end{cases}$ is the Heavyside function.

102 To have the streamfunction of the basic state, we need the Fourier transforms of
103 the buoyancies : for $i = s, b$, we have $\widehat{B}^i(\rho, n) = 0$ for $n \neq 0$ because B^i is independent
104 of ϕ . For $n = 0$, because $\frac{d}{dx}(xJ_1(x)) = xJ_0(x)$, we have $\widehat{B}^i(\rho, 0) = B_0^i \frac{J_1(\rho)}{\rho}$. Then the
105 streamfunction for the basic state at the two boundary levels is :

$$\Psi^s = \int_0^\infty \frac{J_1(\rho)J_0(\rho r)}{\rho\sigma} \left(\frac{B_0^b}{\sinh(\rho\sigma)} - \frac{B_0^s}{\tanh(\rho\sigma)} \right) d\rho \quad (16a)$$

$$\Psi^b = \int_0^\infty \frac{J_1(\rho)J_0(\rho r)}{\rho\sigma} \left(\frac{B_0^b}{\tanh(\rho\sigma)} - \frac{B_0^s}{\sinh(\rho\sigma)} \right) d\rho \quad (16b)$$

106 The flow at the surface of the ocean induced by the vortices is along \vec{e}_ϕ and because
107 $J'_0 = -J_1$, we have :

$$U_r^s = -\frac{1}{r} \partial_\phi \Psi^s = 0 \quad (17a)$$

$$U_\phi^s = \partial_r \Psi^s = \int_0^\infty \frac{J_1(\rho)J_1(\rho r)}{\sigma} \left(\frac{B_0^s}{\tanh(\rho\sigma)} - \frac{B_0^b}{\sinh(\rho\sigma)} \right) d\rho \quad (17b)$$

108 Similarly, the flow at the bottom of the ocean is :

$$U_r^b = 0 \quad (18a)$$

$$U_\phi^b = \int_0^\infty \frac{J_1(\rho)J_1(\rho r)}{\sigma} \left(\frac{B_0^s}{\sinh(\rho\sigma)} - \frac{B_0^b}{\tanh(\rho\sigma)} \right) d\rho \quad (18b)$$

109 We now introduce the quantities :

$$I_n(r, \sigma) := \int_0^\infty \frac{J_n(\rho)J_n(\rho r)}{\sigma \tanh(\rho\sigma)} d\rho \quad (19a)$$

$$M_n(r, \sigma) := \int_0^\infty \frac{J_n(\rho)J_n(\rho r)}{\sigma \sinh(\rho\sigma)} d\rho \quad (19b)$$

110 such that

$$U_\phi^s = B_0^s I_1 - B_0^b M_1 \quad (20a)$$

$$U_\phi^b = B_0^s M_1 - B_0^b I_1 \quad (20b)$$

111 **Remark 3.** This is indeed a steady basic state : since there is no radial velocity, the buoyancy
112 anomaly, which is a tracer, will remain a circular patch if unperturbed.

113 4. Evolution of the vortex boundaries in the linear instability of the vortex

114 Now we perturb each vortex by deforming its contour. Because $\frac{Db^{s,b}}{Dt} = 0$, an initial
 115 plateau in buoyancy will remain such at all times (with a deformed external contour).
 116 Therefore the lateral jump in buoyancy will always exist and we can define the vortex
 117 boundaries as the place where the jump lies. The evolution of the vortex boundaries will
 118 measure the stability of this particular basic state.

119 Assume that the radii R^s and R^b of the vortices at the surface and the bottom are
 120 disturbed from their basic states 1, as represented in Figure 2.

121 **Remark 4.** We assume that during the linear stage of instability, as the perturbation amplitude
 122 remains small, the boundary will not be locally multi-valued, such that we can use the following
 123 parameterisation :

$$\begin{cases} R^s(\phi, t) = 1 + \eta^s(\phi, t) \\ R^b(\phi, t) = 1 + \eta^b(\phi, t) \end{cases} \quad (21)$$

124 where η^i is small compared with 1.

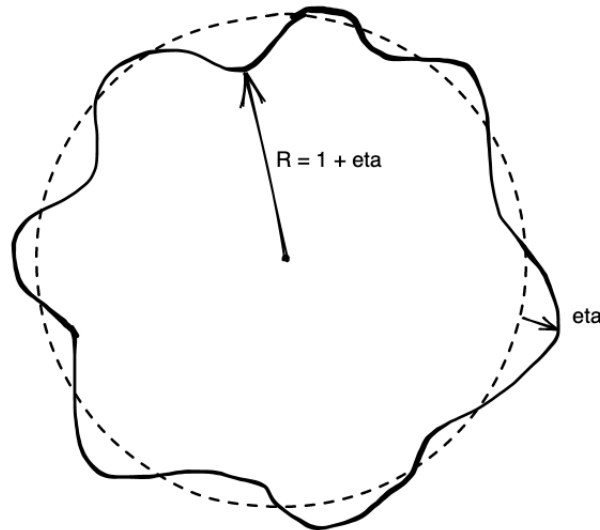


Figure 2. Perturbation of the buoyancy disk, with $\eta \ll 1$.

125 Then the buoyancy at the surface is

$$B_0^s(1 - H(r - R^s(\phi, t))) = \begin{cases} B_0^s & \text{if } r < R^s \\ 0 & \text{if } r > R^s \end{cases} \quad (22)$$

126 where H is the Heavyside function. A similar form can be derived for the buoyancy
 127 at the bottom.

128

129 Because we chose (15) as basic buoyancies, the perturbed buoyancy at the surface is
 130 :

$$b^s = B_0^s(H(r - 1) - H(r - 1 - \eta^s)) = \begin{cases} B_0^s & \text{if } 1 < r < 1 + \eta^s \\ -B_0^s & \text{if } 1 + \eta^s < r < 1 \\ 0 & \text{else} \end{cases} \quad (23)$$

131 As a distribution, for small η^s , we have

$$\begin{cases} b^s(r, \phi, t) = B_0^s \eta^s(\phi, t) \delta_1(r) \\ b^b(r, \phi, t) = B_0^b \eta^b(\phi, t) \delta_1(r) \end{cases} \quad (24)$$

132 where δ_1 is the Dirac distribution in 1.

133 Because the perturbed streamfunctions from Eq. (13) are searched, we need the
134 Fourier transform of the perturbed buoyancies :

$$\widehat{b}^s(\rho, n, t) = \frac{1}{2\pi} \int_0^{2\pi} \int_0^\infty B_0^s \eta^s(\phi, t) \delta_1(r) J_n(\rho r) r dr e^{-in\phi} d\phi \quad (25)$$

$$= \frac{B_0^s}{2\pi} \int_0^{2\pi} J_n(\rho) \eta^s(\phi, t) e^{-in\phi} d\phi \quad (26)$$

$$\widehat{b}^s(\rho, n, t) = B_0^s J_n(\rho) \widehat{\eta}^s(n, t) \quad (27)$$

135 and a similar formula for $\widehat{b}^b(\rho, n, t)$. Because the equations are linear, the perturbed
136 streamfunctions can be computed from the formula (14) :

$$\psi^s(r, \phi, t) = \sum_{n \in \mathbf{N}} \int_0^\infty \frac{J_n(\rho r)}{\sigma \sinh(\rho \sigma)} \left(\widehat{b}^b - \widehat{b}^s \cosh(\rho \sigma) \right) d\rho e^{in\phi} \quad (28)$$

$$\psi^s(r, \phi, t) = \sum_{n \in \mathbf{N}} \left(B_0^b \widehat{\eta}^b M_n - B_0^s \widehat{\eta}^s I_n \right) e^{in\phi} \quad (29)$$

$$\psi^b(r, \phi, t) = \sum_{n \in \mathbf{N}} \left(B_0^b \widehat{\eta}^b I_n - B_0^s \widehat{\eta}^s M_n \right) e^{in\phi} \quad (30)$$

137 where $I_n(r, \sigma)$ and $M_n(r, \sigma)$ are defined in (19).

138

139 The (total and perturbed) radial flow at the boundary of the surface vortex is on the
140 one hand :

$$u_r^s(R^s(\phi, t), \phi, t) = -\frac{1}{1 + \eta^s} \partial_\phi \psi^s(1 + \eta^s, \phi, t) \quad (31)$$

141 and on the other hand

$$u_r^s(R^s(\phi, t), \phi, t) = \frac{DR^s}{Dt} = \frac{\partial R^s}{\partial t} + \frac{1}{R^s} \left(U_\phi^s + u_\phi^s \right) \frac{\partial R^s}{\partial \phi} \quad (32)$$

142 The justification of the first equality $u_r^s(R^s(\phi, t), \phi, t) = \frac{DR^s}{Dt}$ is the following : R^s is the ra-
143 dial coordinates of a material line (the boundary of the vortex). So its rate of Lagrangian
144 displacement in time is exactly the radial velocity of the flow.

145

146 With the equality of Equations (31) and (32), using a Taylor expansion for small
147 amplitude perturbations (details are given in Appendix A), we obtain the following
148 system :

$$\begin{cases} \partial_t \eta^s = -\partial_\phi \psi^s - U_\phi^s \partial_\phi \eta^s \\ \partial_t \eta^b = -\partial_\phi \psi^b - U_\phi^b \partial_\phi \eta^b \end{cases} \quad (33)$$

149 where the functions are applied in $r = 1$.

150

151 Assuming that one mode of perturbation will grow faster than any other, we retain
152 only one Fourier mode $\eta^{s,b}(\phi, t) = \widehat{\eta}^{s,b}(t) e^{in\phi}$ so that we obtain the matrix form with

153 $\widehat{\eta} = \begin{pmatrix} \widehat{\eta}^s \\ \widehat{\eta}^b \end{pmatrix}$:

$$\partial_t \hat{\eta} = in \begin{pmatrix} B_0^s(I_n - I_1) + B_0^b M_1 & -B_0^b M_n \\ B_0^s M_n & B_0^b(I_1 - I_n) - B_0^s M_1 \end{pmatrix} \hat{\eta} \quad (34)$$

154 5. Results

155 5.1. Preliminaries : study of the integrals $I_n - I_1$ and M_n

156 We failed to compute analytically the integrals $I_n - I_1$ and M_n so from now, we do
 157 a numerical study of the stability. Numerically, these integrals have poor convergence
 158 properties. In particular the integral $I_n - I_1$ is not absolutely convergent. We present the
 159 method to compute it in Appendix B. Nevertheless, we obtain a graphical representation
 160 of $I_n - I_1$ and M_n with respect to σ in Figure 3.
 161

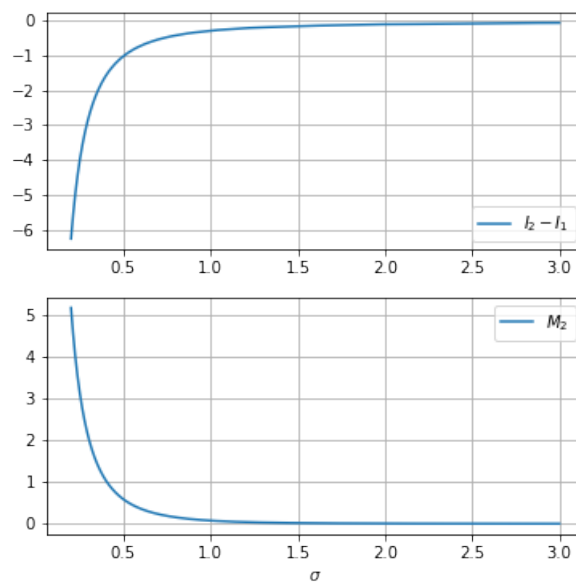


Figure 3. Graphs of the integrals $I_2 - I_1$ and M_2 with respect to σ .

162 5.2. Normal modes

In this section, we consider normal mode perturbations to the vortex boundary. This means that the time dependence of the $\eta^{s,b}$ is $\widehat{\eta}^{s,b}(t) = \mu^{s,b} e^{-i\omega_n t}$ with $\omega_n = a_n + ib_n \in \mathbf{C}$ where b_n is the growth rate and $\mu^{s,b} \in \mathbf{R}^+$. In order to conclude about stability, we are now interested in the sign of b_n , the imaginary part of ω_n . Thanks to (34), we obtain an eigenvalue problem $A_n \mu = -\frac{\omega_n}{n} \mu$ where $\mu = \begin{pmatrix} \mu^s \\ \mu^b \end{pmatrix}$ and

$$A_n = \begin{pmatrix} B_0^s(I_n - I_1) + B_0^b M_1 & -B_0^b M_n \\ B_0^s M_n & B_0^b(I_1 - I_n) - B_0^s M_1 \end{pmatrix}. \quad (35)$$

163 Because $A_n \in \mathcal{M}_n(\mathbf{R})$, there are two possibilities for the eigenvalues. They can be real,
 164 then $b_n = 0$ and the basic state has a neutral stability ; or they can be complex conjugate,
 165 then one of the two eigenvalues has $b_n > 0$ and the basic state is unstable.

Remark 5. The normal mode $n = 1$ is always stable because

$$A_1 = \begin{pmatrix} B_0^b M_1 & -B_0^b M_1 \\ B_0^s M_1 & -B_0^s M_1 \end{pmatrix} \quad (36)$$

166 has two real eigenvalues : 0 and $M_1(B_0^b - B_0^s)$.

Conclusions on this flow stability are obtained by computing the discriminant of $\chi_n(X)$, the characteristic polynomial of A_n :

$$\chi_n(X) = X^2 + (B_0^b - B_0^s)(I_n - I_1 - M_1)X \quad (37)$$

$$- B_0^b B_0^s \left((I_n - I_1)^2 + M_1^2 - M_n^2 \right) - (I_n - I_1) M_1 \left((B_0^s)^2 + (B_0^b)^2 \right). \quad (38)$$

To know if the roots of this polynomial are real or complex conjugate, we compute the sign of the discriminant Δ :

$$\Delta = (B_0^b - B_0^s)^2 (I_n - I_1 - M_1)^2 \quad (39)$$

$$+ 4B_0^b B_0^s \left((I_n - I_1)^2 + M_1^2 - M_n^2 \right) \quad (40)$$

$$+ 4(I_n - I_1) M_1 \left((B_0^s)^2 + (B_0^b)^2 \right) \quad (41)$$

$$\Delta = (I_n - I_1 + M_1)^2 (B_0^b + B_0^s)^2 - 4B_0^b B_0^s M_n^2 \quad (42)$$

167 The conclusion is : if $\Delta > 0$, then the system is neutral because $b_n = 0$ and if $\Delta < 0$,
168 then the growth rate $b_n = \frac{n\sqrt{-\Delta}}{2} \neq 0$ and the system is unstable.

169 **Remark 6.** If we take, as Badin and Poulin in [6] $B_0^b = 0$, we obtain $\Delta = B_0^{s2}(I_n - I_1)^2 > 0$
170 and then we recover their dispersion relation $\omega_n = nB_0^s(I_1 - I_n) \in \mathbf{R}$.

171 Because the case $n = 1$ is always stable, we plot in Figure 4 the normal modes for
172 $n = 2, 3, 4$ and 5.

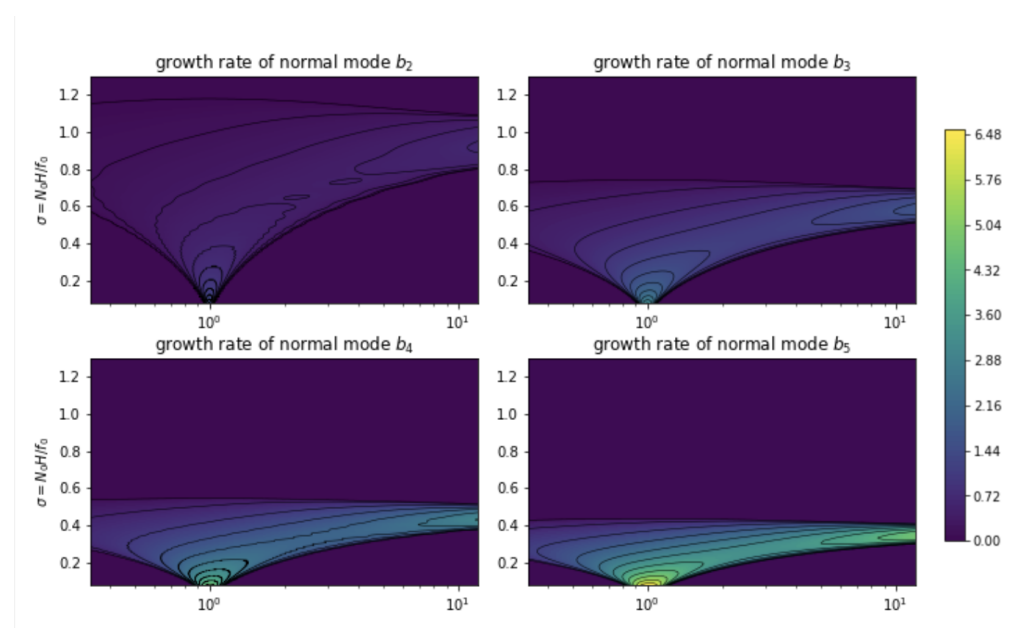


Figure 4. b_n for $n = 2, 3, 4$ and 5, with respect to σ and $\frac{B_0^s}{B_0^b}$.

The dark purple zone is where $b_n = 0$. There, $\Delta > 0$ and the system reaches a stable state with the following dispersion relation :

$$\omega_n(\sigma) = \frac{n}{2} \left[\left(B_0^s - B_0^b \right) (I_n - I_1 - M_1) \pm \sqrt{\Delta} \right] \quad (43)$$

173 The four angular modes have similar stability properties on the top-hat
174 vortex, but for different values of the physical parameters $(\sigma, B_0^s/B_0^b)$. For each mode, we
175 can separate three stable zones :

- 176 • in the right hand side of each panel, where σ is larger than a threshold σ_{critic}
177 depending on n . We recover here the results of [6] or [7]. They found that a top-hat
178 vortex, alone in a SQG model, is stable. In this area, the system is linearly stable for
179 barotropic (horizontal shear) instability. They are sufficiently far from the other (σ is
180 proportional to H) so we could neglect the interactions. The two-layer SQG model
181 is then viewed as two one-layer SQG models where there are two independent
182 top-hat vortices. We define σ_{critic} as the critical value of σ leading to instability,
183 all other parameters being fixed. σ_{critic} is a decreasing function with respect to n .
184 The stability of high mode perturbations is reached for a smaller distance between
185 vortices than low mode perturbations. This is due to the relation between horizontal
186 and vertical wave numbers in the SQG model.
- 187 • in the top left and the bottom left sides of each panel, the system is also stable. In
188 these areas, the vortices are close to each other but have very different intensities.
189 An interpretation of this situation could be that perturbations on one of the vortice
190 have very different phase speeds around the contour than for the other vortex. The
191 impossibility for these two (Rossby) waves to phase lock and resonate stabilizes the
192 whole system.

193 The system is unstable if the mean buoyancy intensities are similar and if the
194 vortices are vertically close to each other.

195 For a given fixed mode and a fixed ratio $B_0^s/B_0^b > 1$, an interpretation of the change
196 of stability is the following :

- 197 • for small σ , the two vortices are too close to each other for the wave to grow; thus
198 short wave cut off (usual for the Eady model) can be explained by the absence of
199 phase locking between waves.
- 200 • for intermediate σ , the distance between the two vortices allows the mode to grow
201 (phase locking with the proper phase shift is possible) and then the system is
202 unstable. The smaller σ is, the shorter the most unstable waves are.
- 203 • for large σ , the two vortices are far from each other, wave-wave interaction is weak
204 and the mode is stable.

205 5.3. Singular modes

206 System (34) is a 2×2 system, and the matrix A_n is independent of time t . The
207 solution is then given by $\widehat{eta} = e^{A_n t} \widehat{eta}_0$. Since matrix A is not self-adjoint, linear
208 combinations of normal modes can grow faster than the normal modes [13]; they are the
209 singular modes of the problem. Therefore we calculate the singular modes of stability
210 of this vortex flow. They are defined as the maximal growth rate of a given norm of
211 these perturbations. Here we choose the squared perturbation buoyancy as the norm.
212 The numerical method for singular mode calculation given in the appendix of [14] is
213 implemented here. The grow rates of the singular modes are the eigenvalues of the
214 matrix $e^{A_n t}$. These growth rates are shown with respect to the same physical parameters
215 $(\sigma, B_0^s/B_0^b)$ as the normal modes, for increasing values of the time t and for $n = 2$ on
216 figure 5.

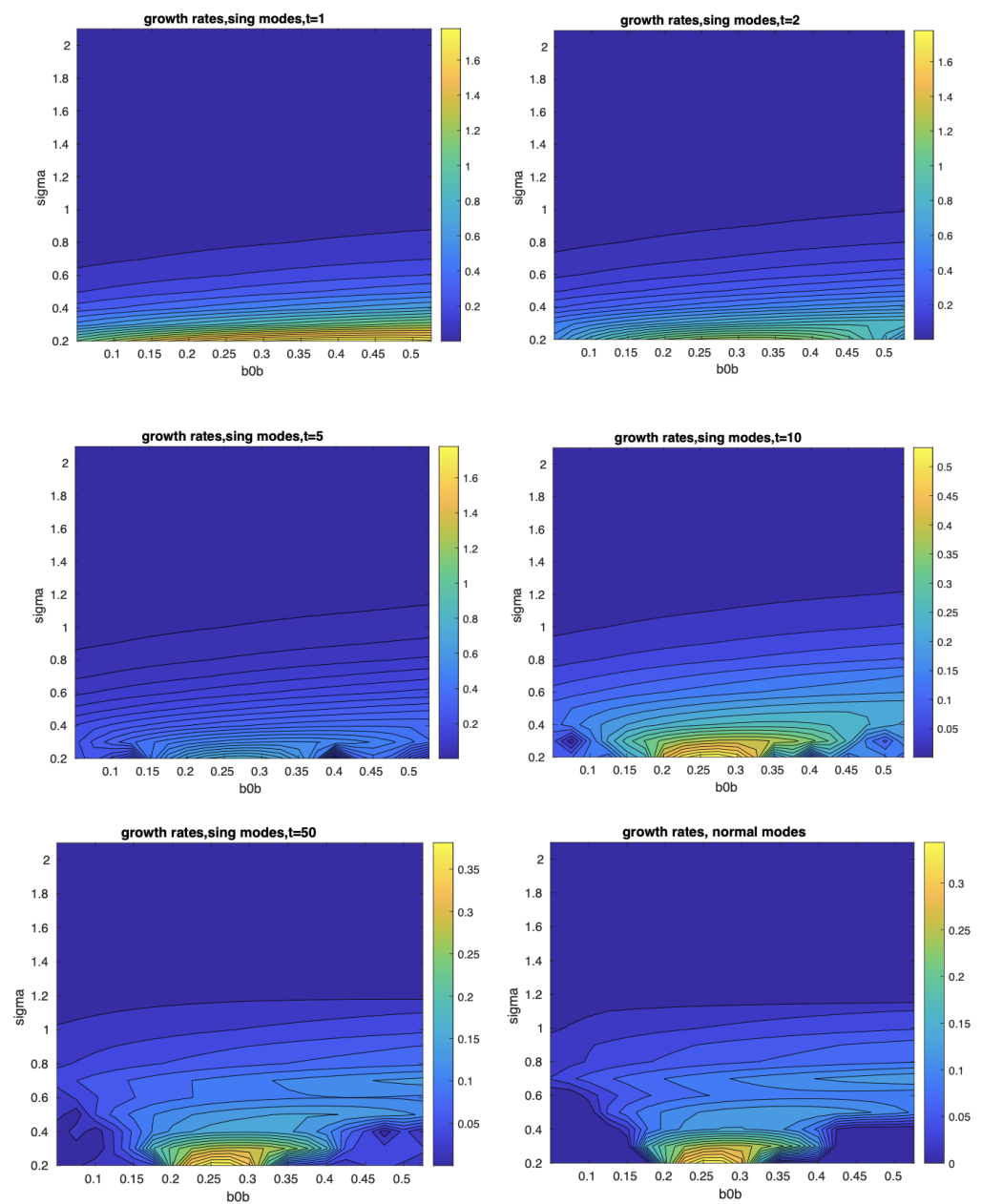


Figure 5. Singular mode for $n = 2$ for different times $t = 1, 2, 5, 10, 50$. Here $B_0^s = 0.25$. The bottom right panel represents the growth rates for normal mode $n = 2$. We can note the convergence of singular modes to normal modes.

217 For short time, the singular modes growth rates are concentrated near the region of
 218 small σ . Indeed, at small times, only short waves can grow corresponding to buoyancy
 219 surfaces close to each other vertically. Also, for small σ , the singular mode instability
 220 occurs for every B_0^b , i.e. even a weak bottom buoyancy is sufficient then to allow
 221 the phase locking of these short waves. Also this implies that two vertically close
 222 vortices, with different mean buoyancies, are unstable for singular modes but stable in
 223 normal modes for small time. Note that a similar remark was made in [15] about the
 224 independence of singular mode growth rates to the barotropic component of the flow at
 225 short times. As time grows, the singular mode growth rates for $n = 2$ converge towards
 226 those of the corresponding normal modes; this result is valid for the other $n \in \mathbb{N}$. This
 227 confirms the result of previous studies [13–15].

228 6. Conclusions and perspectives

229 In this study, we have developed analytically and numerically the calculation
 230 of growth rates for the instability of two superposed vortices. The theory and the
 231 computation of the mean flow were first done analytically. The computation of the
 232 normal and singular modes was done numerically. The two-level SQG model and
 233 the considered steady state are idealized, but they provide simple results for the Eady
 234 baroclinic instability of two superimposed vortices : stability for vertically distant
 235 vortices, instability for vertically close vortices, similar in intensity and instability in
 236 singular modes only for small time for close and different in intensity vortices.

237 Though these results pertain to an idealized vortex, we can apply them to the ocean.
 238 Using the following values $f_0 = 10^{-4} s^{-1}$, $N_0 = 5 \cdot 10^{-3} s^{-1}$, $R = 2.5 \cdot 10^4 m$, $h = 10^3 m$, $V =$
 239 $0.5 m/s$, where V is the rotational velocity of an oceanic vortex, we obtain the following
 240 length and time scales, $L = 2.5 \cdot 10^4 m$, $T = 6.3 \cdot 10^4 s$. Firstly, we can use the values of
 241 σ to determine which vortices can be unstable: strong instability occurs for $\sigma = 0.2$
 242 leading to most unstable vortices having a thickness of 100 m which indeed corresponds
 243 to small vortices (with radii close to 10 km). Secondly, we can compute the growth rates
 244 of such normal mode perturbations in the ocean: in dimensionless terms, they are on the
 245 order of 0.3 for $n = 2$. This corresponds to a typical time scale for the growth of these
 246 perturbations of $6.3/0.3 \cdot 10^4 s$, about 2.5 days. This timescale is slightly shorter than that
 247 found in the two-layer Phillips problem of mesoscale vortex baroclinic instability, about
 248 4 days [1,15].

249 Finally, we must also note that we have studied only the linear instability of such
 250 vortices. The natural follow-up of these calculations is the study of the long-term,
 251 nonlinear evolution of these unstable vortices. This will indicate if the linearly unstable
 252 waves found here, can be stabilised on the long run via nonlinear wave interactions and
 253 which shape the nonlinearly stabilised vortices would take. Furthermore, this will allow
 254 the variation of several parameters not included here:

- 255 • investigate the effect of different radii for the two vortices
- 256 • shift one vortex with respect to the other and study the evolution of tilted vortices
- 257 • consider two different modes n_s and n_b of perturbation for the two vortices
- 258 • consider other radial shapes for the vortices (Gaussian,...)

259 This numerical study is under way and will complement this first paper.

260 **Funding:** This research has been supported by external funding: the first author has received a
 261 Ph.D grant from Ecole Normale Supérieure, Rennes

262 **Conflicts of Interest:** The authors declare no conflict of interest. The funders had no role in the
 263 design of the study; in the collection, analyses, or interpretation of data; in the writing of the
 264 manuscript, or in the decision to publish the results.

265 Abbreviations

266 The following abbreviation is used in this manuscript:

267
 268 SQG Surface Quasi Geostrophic

269 Appendix A. Proof of the system (33)

We will develop the calculus for the surface only. The bottom case is similar. On
 one hand, from Equation (31), we have :

$$u_r^s(R^s(\phi, t), \phi, t) = -(1 + O(\varepsilon)) \left(\partial_\phi \psi^s(1, \phi, t) + O(\varepsilon^2) \right) \quad (A1)$$

$$= -\partial_\phi \psi^s(1, \phi, t) + O(\varepsilon^2) \quad (A2)$$

On the other hand, from Equation (32) we have :

$$u_r^s(R^s(\phi, t), \phi, t) = \frac{\partial \eta^s}{\partial t}(\phi, t) + \frac{1}{1 + \eta^s(\phi, t)} \quad (\text{A3})$$

$$\left[U_\phi^s(1 + \eta^s(\phi, t), \phi, t) + u_\phi^s(1 + \eta^s(\phi, t), \phi, t) \right] \frac{\partial \eta^s}{\partial \phi}(\phi, t) \quad (\text{A4})$$

$$= \partial_t \eta^s(\phi, t) + (1 + O(\varepsilon)) \left[U_\phi^s(1, \phi, t) + O(\varepsilon) \right] \partial_\phi \eta^s(\phi, t) \quad (\text{A5})$$

$$= \partial_t \eta^s(\phi, t) + U_\phi^s(1, \phi, t) \partial_\phi \eta^s(\phi, t) + O(\varepsilon^2) \quad (\text{A6})$$

And then, at the order $O(\varepsilon)$, we obtain

$$\partial_t \eta^s(\phi, t) = -\partial_\phi \psi^s(1, \phi, t) - U_\phi^s(1, \phi, t) \partial_\phi \eta^s(\phi, t) \quad (\text{A7})$$

270 **Remark A1.** Note a difficulty we did not mention earlier : U_ϕ^s is not differentiable in the classical
271 way. Formally, we should work with smooth approximation of top-hat vortices and then move on
272 to the limit.

273 **Appendix B. Proof of convergence and numerical method to compute $I_n - I_1$ and M_n**

274 Recall $I_n - I_1 = \frac{1}{\sigma} \int_0^\infty f_n(x) dx$ and $M_n = \frac{1}{\sigma} \int_0^\infty g_n(x) dx$ with $f_n(x) = \frac{J_n(x)^2 - J_1(x)^2}{\tanh(\sigma x)}$
275 and $g_n(x) = \frac{J_n(x)^2}{\sinh(\sigma x)}$ (see Figure A1). The two functions are continuous on $]0, \infty[$. Let's
276 do the analysis of convergence in 0 and in ∞ :

- For x in a neighborhood of 0, $J_n(x) \sim \frac{x^n}{2^{n-1} n!}$ so for $n > 1$:

$$f_n(x) \underset{0}{\sim} -\frac{J_1(x)^2}{\tanh(\sigma x)} \quad (\text{A8})$$

$$f_n(x) \underset{0}{\sim} -\frac{x}{4\sigma} \quad (\text{A9})$$

and for $n \geq 1$:

$$g_n(x) \underset{0}{\sim} \frac{x^{2n-1}}{4^n (n!)^2 \sigma}. \quad (\text{A10})$$

277 So the two functions are integrable in 0.

- For x in a neighborhood of $+\infty$,

$$J_n(x) \underset{\infty}{=} \sqrt{\frac{2}{\pi x}} \sin\left(x - \frac{n\pi}{2} + \frac{\pi}{4}\right) - \frac{4n^2 - 1}{4\sqrt{2\pi x^{\frac{3}{2}}}} \sin\left(x - \frac{n\pi}{2} - \frac{\pi}{4}\right) + o\left(\frac{1}{x^{\frac{3}{2}}}\right) \quad (\text{A11})$$

The computation gives for $n = 2p > 1$:

$$f_{2p}(x) \underset{\infty}{\sim} \frac{2 \sin(2x)}{\pi x} \quad (\text{A12})$$

and for $n = 2p + 1 \geq 1$:

$$f_{2p+1}(x) \underset{\infty}{\sim} \frac{1 - (2p + 1)^2}{\pi x^2} \cos(2x). \quad (\text{A13})$$

We can quickly conclude for the odd case because $f_{2p+1} = O\left(\frac{1}{x^2}\right)$ is absolutely convergent in $+\infty$. The even case is a modified integral sine so that it converges. For every $n \in \mathbb{N}$, we have a quick convergence in $+\infty$ for g_n :

$$g_n(x) \underset{\infty}{\sim} \frac{4}{\pi x} \sin^2\left(x - \frac{n\pi}{2} + \frac{\pi}{4}\right) e^{-\sigma x}. \quad (\text{A14})$$

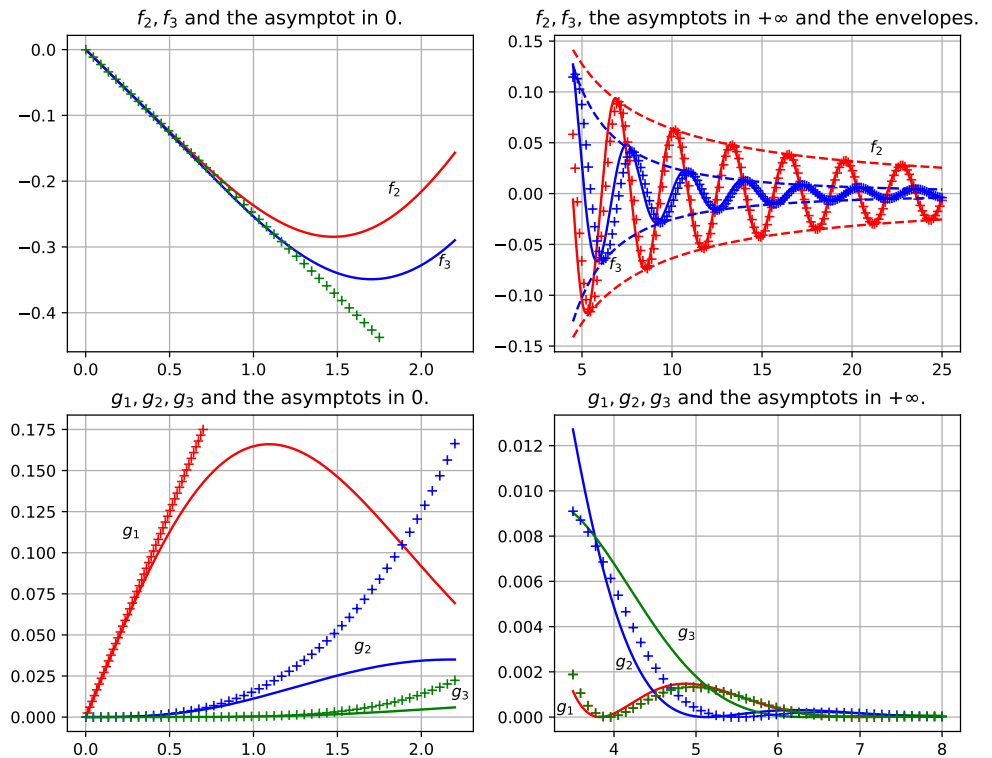


Figure A1. The integrands f_n and g_n are in solid line, the asymptotes are plotted with crosses and the envelopes for the top right panel are in dashed lines. Notice that for the bottom right panel, the asymptote depends only on the parity of n . This explains why there are only two asymptotes plotted. Here we take the parameter $\sigma = 1$.

278 Numerically, the only difficult point is the fact that the integral sine does not
 279 absolutely converge so that the classically implemented methods to compute integrals
 280 are not adapted. A python routine exists to compute integral sine and this is what we use.
 281 The idea is to cut the integrals in three parts, $\int_0^\infty = \int_0^\epsilon + \int_\epsilon^A + \int_A^\infty$, to use approximation
 282 for the integrals in 0 and in $+\infty$ and to use classical python routine in $[\epsilon, A]$. With the
 283 asymptotic developments we used, we obtain :

- for f_n in 0 :

$$\int_0^\epsilon f_n(x) dx \simeq -\frac{\epsilon^2}{8\sigma} \quad (\text{A15})$$

- for f_{2p} in $+\infty$:

$$\int_A^\infty f_{2p}(x) dx \simeq \frac{2}{\pi} \int_{2A}^\infty \frac{\sin t}{t} dt \quad (\text{A16})$$

$$\simeq 1 - \frac{2}{\pi} \text{Si}(2A) \quad (\text{A17})$$

284 where $\text{Si}(x) = \int_0^x \frac{\sin(t)}{t} dt$ is the integral sine function.

- for f_{2p+1} in $+\infty$:

$$\int_A^\infty f_{2p+1}(x)dx \simeq \frac{1 - (2p+1)^2}{\pi} \left(- \int_A^\infty \frac{2 \sin(2x)}{x} dx - \left[\frac{\cos 2x}{x} \right]_A^\infty \right) \quad (\text{A18})$$

$$\simeq \frac{1 - (2p+1)^2}{\pi} \left(2 \text{Si}(2A) - \pi + \frac{\cos 2A}{A} \right) \quad (\text{A19})$$

- for g_n in 0 :

$$\int_0^\varepsilon g_n(x)dx \simeq \frac{\varepsilon^{2n}}{2n(n!)^2 4^n \sigma} \quad (\text{A20})$$

- for g_n in $+\infty$:

$$\int_A^\infty g_n(x)dx \simeq \frac{4}{\pi} \int_A^\infty \frac{\sin^2 \left(x - \frac{n\pi}{2} + \frac{\pi}{4} \right)}{x} e^{-\sigma x} dx \quad (\text{A21})$$

$$\leq \frac{4}{\pi} \int_A^\infty \frac{1}{x} e^{-\sigma x} dx \quad (\text{A22})$$

$$\leq \frac{4 e^{-\sigma A}}{\pi \sigma A} \quad (\text{A23})$$

285 So if we take σA sufficiently large (in practise we take $\sigma A \simeq 20$), this part can be
286 neglected.

287 The following Table 1 sums up the approximations we used to compute the integrals
288 $I_n - I_1$ and M_n .

	f_{2p}	f_{2p+1}	g_n
\int_0^ε	$-\frac{\varepsilon^2}{8\sigma}$	$-\frac{\varepsilon^2}{8\sigma}$	$\frac{\varepsilon^{2n}}{2n(n!)^2 4^n \sigma}$
\int_A^∞	$1 - \frac{2}{\pi} \text{Si}(2A)$	$\frac{1 - (2p+1)^2}{\pi} \left(2 \text{Si}(2A) - \pi + \frac{\cos 2A}{A} \right)$	0

Table 1: Summary of the approximated integrals

289

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