

Mathematics of Waves and Fields

Dominik Szablonski

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Chapter 1

Partial Differential Equations

1.1 Seperation of Variables

Suppose we have a PDE whose solution is in the form, $u(r_1, r_2, \dots, r_n)$ where there are n co-ordinates r_i , then we can solve the PDE by seperation of variables by assuming a solution of the form,

$$u(r_1, r_2, \dots, r_n) = R_1(r_1)R_2(r_2) \cdots R_n(r_n). \quad (1.1)$$

This will turn a compatible PDE into an ODE.

1.1.1 Specific solutions

We are often most interested in the specific solutions to a wave equation. In order to get a specific solution, constraints/boundary conditions must be provided. The general method is as follows,

1. Use seperation of variables;
2. Build superpositions of solutions;
3. Apply boundary conditions and find appropriate constants.

1.2 Series Solutions

The general steps to solving an ODE using this method are,

1. Assume a series solution of the form,

$$y = \sum_{n=0}^{\infty} a_n x^n \quad (1.2)$$

2. Obtain the recurrence relation.

Chapter 2

Fourier Series

Given a periodic function $f(x)$ with period $2L$ in the range $-L \leq x \leq L$, the fourier expansion is given by,

$$f(x) = \frac{a_0}{2} + \sum_{n=1}^{\infty} \left(a_n \cos \frac{n\pi x}{L} + b_n \sin \frac{n\pi x}{L} \right) \quad (2.1)$$

for,

$$a_n = \frac{1}{L} \int_{-L}^L f(x) \cos \frac{n\pi x}{L} dx \quad b_n = \frac{1}{L} \int_{-L}^L f(x) \sin \frac{n\pi x}{L} dx \quad (2.2)$$

In order to expand a function it must meet *Dirichlet's Conditions*, so the function must,

1. be single valued,
2. have a finite number of discontinuities,
3. $\int_{-L}^L |f(x)| dx$ must be finite.

We say that $\sin \frac{n\pi x}{L}$ and $\cos \frac{n\pi x}{L}$ form a *complete, orthogonal basis*. And furthermore, the Fourier series allows *an expansion of a function on a set of orthogonal basis functions*.

2.1 Exponential Fourier Series

We can further write the Fourier expansion in terms of complex exponentials,

$$f(x) = \sum_{n=-\infty}^{\infty} c_n \exp \left(i \frac{n\pi x}{L} \right) \quad (2.3)$$

for,

$$c_n = \frac{1}{2L} \int_{-L}^L f(x) \exp \left(-i \frac{n\pi x}{L} \right) dx \quad (2.4)$$

From our complex definition of the Fourier series, we can say that 2 complex functions $u(z)$ and $v(z)$ are orthogonal on the interval $a \leq z \leq b$ if,

$$\int_a^b u(z)v(z) dz = 0. \quad (2.5)$$

Chapter 3

Fourier Transform

If we wish to analyse non-periodic functions, we can take the limit of our range, $\lim_{L \rightarrow \infty}(-L, L)$. Let us write,

$$k_n = \frac{n\pi}{L} \qquad \Delta k = \frac{\pi}{L} \qquad (3.1)$$

so,

$$(3.2)$$

$$c_n = \frac{1}{2L} \int_{-L}^L f(x) \exp(-ik_n x) dx. \qquad (3.3)$$

Let us write $F(k) = 2Lac_n$,

$$F(k_n) = a \int_{-L}^L f(x) \exp(-ik_n x) dx \qquad (3.4)$$

$$\implies f(x) = \frac{1}{2\pi a} \sum_{n=-\infty}^{\infty} F(k_n) \exp(ik_n x) \Delta k. \qquad (3.5)$$

In the limit of $L \rightarrow \infty$, we obtain,

$$F(k) = a \int_{-\infty}^{\infty} f(x) \exp(-ikx) dx$$

Fourier transform of $f(x)$. (3.6)

$$f(x) = \frac{1}{2\pi a} \int_{-\infty}^{\infty} F(k) \exp(ikx) dk$$

Inverse fourier transform of $F(k)$. (3.7)

We define constant,

$$a = \begin{cases} \text{unity} & \text{Physics} \\ \frac{1}{\sqrt{2\pi}} & \text{Maths} \end{cases}. \qquad (3.8)$$

3.1 Fourier Transform Theorems

3.1.1 Similarity Theorem

Theorem.

$$\mathcal{F}\{f(ax)\} = \frac{1}{a} F\left(\frac{k}{a}\right).$$

(3.9)

3.1.2 Reverse fourier transform over all space

Theorem.

$$\int_{-\infty}^{\infty} f(x) dx = F(0)$$

(3.10)

3.1.3 Translation of fourier transform

Theorem.

$$\boxed{\mathcal{F}\{f(x-a)\} = e^{-ika}F(k)} \quad (3.11)$$

3.1.4 General manipulation of the fourier transform

Most generally,

$$\boxed{\mathcal{F}\{f(ax-b)\} = e^{-ikb}\frac{1}{a}F\left(\frac{k}{a}\right)} \quad (3.12)$$

3.2 Dirac-Delta Function

Dirac's original approximation of the δ function used a function $\Pi(x)$ which was defined,

$$\Pi(x) = \begin{cases} 1 & -\frac{1}{2} \leq x \leq \frac{1}{2} \\ 0 & \text{Elsewhere} \end{cases} \quad (3.13)$$

Using this definition we can write,

$$\delta(x) = \lim_{k \rightarrow \infty} \{k\Pi(kx)\} \quad (3.14)$$

It can also be defined using sinc,

$$\delta(x) = \lim_{k \rightarrow \infty} \left\{ \frac{k}{\pi} \frac{\sin kx}{kx} \right\} \quad (3.15)$$

However, the most commonly used, and most applicable form is,

$$\boxed{\delta(x) = \frac{1}{2\pi} \int_{-\infty}^{\infty} e^{ikx} dx} \quad (3.16)$$

We note 4 important properties of the Dirac-Delta,

1. $\lim_{x \rightarrow 0} \delta(x) = \infty$
2. $\int_{-\infty}^{\infty} \delta(x) dx = 1$
3. $\delta(ax+b) = \frac{1}{|a|} \delta\left(x + \frac{a}{b}\right)$
4. $\delta(x) = \delta(-x)$ i.e., Dirac-Delta is even.

3.3 Parseval's Theorem

Theorem.

$$\boxed{\int_{-\infty}^{\infty} |f(x)|^2 dx = a \int_{-\infty}^{\infty} |F(k)|^2 dk} \quad (3.17)$$

where $a = 1$ for mathematical symmetry, and $a = \frac{1}{2\pi}$ for physical symmetry.

Proof.

$$\begin{aligned} \int_{-\infty}^{\infty} |f(x)|^2 dx &= \frac{1}{2\pi} \int_{-\infty}^{\infty} \int_{-\infty}^{\infty} F(k) e^{ikx} dk \int_{-\infty}^{\infty} F^*(k') e^{ik'x} dk' dx \\ &= \int_{-\infty}^{\infty} F(k) dk \int_{-\infty}^{\infty} F^*(k') dk' \underbrace{\frac{1}{2\pi} \int_{-\infty}^{\infty} e^{i(k-k')x} dx}_{\delta(k-k')} \\ &= \int_{-\infty}^{\infty} |F(k)|^2 dk \end{aligned} \quad (3.18)$$

□

3.4 Convolution

We define the convolution $h(x)$ of two functions $f(x)$ and $g(x)$ as,

$$\begin{aligned} h(x) &= f(x) * g(x) = \int_{-\infty}^{\infty} f(x - x')g(x') \, dx' \\ &= \int_{-\infty}^{\infty} f(x')g(x - x') \, dx'. \end{aligned} \quad (3.19)$$

If we define the fourier transforms of $f(x)$ and $g(x)$,

$$F(k) = \mathcal{F}\{f(x)\} \quad G(k) = \mathcal{F}\{g(x)\} \quad (3.20)$$

then the Fourier transform of the convolution is given by,

$$\mathcal{F}\{h(x)\} = \mathcal{F}\{f(x) * g(x)\} = F(k)G(k). \quad (3.21)$$

Proof. Let us define $\zeta = x - x'$, then,

$$\begin{aligned} H(k) &= \int_{-\infty}^{\infty} \int_{-\infty}^{\infty} dx \, dx' f(\zeta)g(x')e^{-ikx} \\ &= \int_{-\infty}^{\infty} \int_{-\infty}^{\infty} dx' \, d\zeta f(\zeta)g(x')e^{-ik(\zeta+x')} \\ &= \int_{-\infty}^{\infty} g(x')e^{ikx'} \, dx' \int_{-\infty}^{\infty} f(\zeta)e^{ik\zeta} \, d\zeta \end{aligned} \quad (3.22)$$

which clearly corresponds to the product of the two transforms. \square

3.4.1 Example: Diffraction through 2 slits

3.5 Wave Packets

In 1 dimension, a forward travelling wave is defined by,

$$\phi(x, t) = e^{-i(kx - \omega t)} \quad (3.23)$$

which satisfies the 1 dimensional wave equation,

$$\frac{\partial^2 \phi}{\partial x^2} = \frac{1}{c^2} \frac{\partial^2 \phi}{\partial t^2}. \quad (3.24)$$

By substituting in the travelling wave, we find,

$$k^2 = \frac{1}{c^2} \omega^2 \implies \omega = ck. \quad (3.25)$$

A plane wave $\phi(\mathbf{x}, t) = e^{i(\mathbf{k} \cdot \mathbf{x} - \omega t)}$ satisfies the 3 dimensional wave equation,

$$\nabla^2 \phi = \frac{1}{c^2} \frac{\partial^2 \phi}{\partial t^2}. \quad (3.26)$$

Where we have,

$$\omega = c|k| \quad (3.27)$$

for a plane travelling along k .

Returning to the 1 dimensional wave, we can sum these travelling waves along the $+x$ direction,

$$\phi(x, t) = \frac{1}{\sqrt{2\pi}} \int_{-\infty}^{\infty} G(k)e^{ik(x-ct)} \, dk \quad (3.28)$$

where $G(k)$ is the Fourier transform of $\phi(x, 0)$. This wave satisfies the wave equation as all components of the wave travel at the same velocity c . We are also able to use the wave equation to describe waves in *non-dispersive media*, i.e., those where the velocity of the waves depends on wavelength,

$$v_p(k) = \frac{\omega(k)}{k}. \quad (3.29)$$

The most general wave can be written as,

$$\phi(x, t) = \frac{1}{\sqrt{2\pi}} \int_{-\infty}^{\infty} G(k) e^{i(kx - \omega(k)t)} dk. \quad (3.30)$$

3.5.1 Dispersion

A dispersive wave packet will have the following properties,

- The envelope wave of the wavepacket will move with group velocity,

$$v_g = \frac{d\omega}{dk} = v_p + k \frac{dv_p}{dk}. \quad (3.31)$$

- The dispersive effects of the wave are a second order effect. i.e., we must expand any approximations to the second order. We will always assume $\omega \equiv \omega(k)$.

Chapter 4

Special Functions

4.1 Taylor Expansion

The Taylor expansion about x_0 is,

$$\begin{aligned} f(x) &= f(x_0) + (x - x_0)f'(x_0) + \frac{(x - x_0)^2}{2!} + \dots \\ &= \sum_{n=0}^{\infty} f^{(n)}(x_0) \frac{(x - x_0)^n}{n!} \end{aligned} \quad (4.1)$$

Let us then redefine this as a simple series,

$$f(x) = \sum_{n=0}^{\infty} u_n \quad (4.2)$$

we can then define the convergence criteria,

$$\lim_{n \rightarrow \infty} |r_n| = \lim_{n \rightarrow \infty} \left| \frac{u_{n+1}}{u_n} \right| < 1. \quad (4.3)$$

4.2 Hermit's Equation

$$\frac{d^2 y}{dx^2} - 2x \frac{dy}{dx} + 2ny = 0, \forall y, n \in \mathbb{Z} \quad (4.4)$$

We can obtain solutions to Hermit's equation by assuming a series solution,

$$y = \sum_{k=0}^{\infty} x^k. \quad (4.5)$$

Substituting this into Hermit's equation,

$$\sum_{k=0}^{\infty} k(k-1)a_k x^{k-2} - 2x \sum_{k=0}^{\infty} k a_k x^{k-1} + 2n \sum_{k=0}^{\infty} a_k x^k = 0 \quad (4.6)$$

Let us shift k such that, $k \rightarrow k+2$,

$$\begin{aligned} \sum_{k=0}^{\infty} (k+2)(k+1)a_{k+1} x^k - 2 \sum_{k=0}^{\infty} k a_k x^k + 2n \sum_{k=0}^{\infty} a_k x^k &= 0 \\ \sum_{k=0}^{\infty} \{(k+2)(k+1)a_{k+2} - (2k-2n)a_k\} x^k &= 0 \end{aligned} \quad (4.7)$$

from which we obtain the recurrence relation,

$$a_{k+2} = \frac{2(k-n)}{(k+2)(k+1)} a_n. \quad (4.8)$$

Using this recurrence relation, we are able to form a solution for y , starting at $k = 0$ and $k = 1$ to obtain the even and odd solutions respectively,

$$y = a_0 \left[1 - \frac{2n}{2!}x^2 - \frac{2n(4-2n)}{4!}x^4 + \dots \right] + a_1 \left[x + \frac{(2-2n)}{3!}x^3 + \frac{(2-2n)(6-2n)}{5!}x^5 + \dots \right]. \quad (4.9)$$

Let us note that at $k = n$ the series will terminate.

Given the solution to Hermit's equation, by considering different values of n , we are able to obtain *Hermit's Polynomials* which are discussed in the section below.

4.2.1 Hermit's Polynomials

We denote Hermit's polynomials by $y \equiv H_n(x)$. By simply looking at eq. (4.9), we see that the first three even Hermit polynomials are,

$$H_0(x) = 1 \quad H_2(x) = 1 - 2x^2 \quad H_4(x) = 1 - 4x^2 + \frac{4}{3}x^4, \quad (4.10)$$

and the first 3 odd ones are,

$$H_1(x) = x \quad H_3(x) = x - \frac{2}{3}x^3 \quad H_5(x) = x - \frac{4}{3}x^3 + \frac{4}{5}x^5 \quad (4.11)$$

In physics, we often normalise the Hermit polynomials such that the highest order term is positive and has a coefficient 2^n .

Orthogonality of Hermit Polynomials

Hermit Polynomials satisfy the orthogonality relation,

$$\int_{-\infty}^{\infty} H_n(x) H_m(x) e^{x^2} dx = \sqrt{\pi} n! 2^m \delta_{nm}. \quad (4.12)$$

This means that Hermit's polynomials can be used as a basis for series expansion of a function. We can further define a normalised Hermit function,

$$\psi_m(x) = \left(\frac{1}{\sqrt{\pi} m! 2^m} \right)^{1/2} H_m e^{\frac{x^2}{2}} \quad (4.13)$$

which satisfies,

$$\int_{-\infty}^{\infty} \psi_m(x) \psi_n(x) dx = \delta_{mn} \quad (4.14)$$

4.3 Legendre's Equation

$$\boxed{(1-x^2) \frac{d^2 y}{dx^2} - 2x \frac{dy}{dx} + \ell(\ell+1)y = 0 \quad \ell \geq 0, \ell \in \mathbb{Z}} \quad (4.15)$$

We solve Legendre's equation by series expansion, from which we obtain,

$$(1-x^2) \sum_{n=2} n(n-1) a_n x^{n-2} - 2x \sum_{n=1} n a_n x^{n-1} \ell(\ell+1) \sum_{n=0} a_n x^n. \quad (4.16)$$

Let us shift the sums so we only have terms in powers of n ,

$$\sum_{n=0} \{(n+2)(n+1)a_{n+2} - n(n-1)a_n - 2na_n x^n - 2na_n x^n + \ell(\ell+1)a_n\} x^n = 0. \quad (4.17)$$

From which we can easily obtain a general recurrence relation,

$$a_{n+2} = \frac{n(n+1) - \ell(\ell+1)}{(n+2)(n+1)} a_n \quad (4.18)$$

thus the solution for Legendre's equation has even and odd parts given below,

$$\begin{aligned} y = & a_0 \left[(1 - \ell(\ell+1)) \frac{x^2}{2!} + (\ell-2)(\ell(\ell+1)(\ell+3)) \frac{x^4}{4!} + \dots \right] & \text{Even} \\ & + a_1 \left[x - (\ell-1)(\ell+2) \frac{x^3}{3!} + (\ell-3)(\ell-1)(\ell+2) \frac{x^5}{5!} + \dots \right] & \text{Odd} \end{aligned} \quad (4.19)$$

which we can clearly see terminates at $\ell = n$, which allows the series to converge.

4.3.1 Legendre Polynomials

The steps to finding Legendre polynomials $P_n(x)$ are as follows,

- Decide whether the polynomial is odd or even, and choose which part of y you will use.
- Find the coefficients of the polynomial $y(x)$ in terms of a_0 for even polynomials and a_1 for odd polynomials.
- Set $y(0) = 1$ to find a value for a_1 or a_0 .
- Evaluate the final polynomial.

Orthogonality of Legendre Polynomials

Legendre polynomials are orthogonal over the interval $|x| \leq 1$, i.e.,

$$\int_{-1}^1 P_l(x) P_m(x) dx = 0 \quad m \neq l. \quad (4.20)$$

Let us recall eq. (4.15), and rewrite,

$$\frac{d}{dx} \left[(1-x^2) \frac{\partial P_l}{\partial x} \right] = -l(l+1) P_l(x) \quad (4.21)$$

$$\frac{d}{dx} \left[(1-x^2) \frac{\partial P_m}{\partial x} \right] = -m(m+1) P_m(x). \quad (4.22)$$

Multiply eq. (4.21) by P_m , and eq. (4.22) by P_l , and take them away from each other. We have,

$$\begin{aligned} \text{LHS} = & \int_{-1}^1 \frac{d}{dx} \left[(1-x^2) \frac{\partial P_l}{\partial x} \right] P_m dx \\ & - \int_{-1}^1 \frac{d}{dx} \left[(1-x^2) \frac{\partial P_m}{\partial x} \right] P_l dx \end{aligned} \quad (4.23)$$

Evaluating eq. (4.23) by parts, we have,

$$u = P_m \quad \frac{dv}{dx} = \frac{d}{dx} \left[(1-x^2) \frac{\partial P_l}{\partial x} \right] \quad (4.24)$$

$$\frac{du}{dx} = \frac{\partial P_m}{\partial x} \quad v = (1-x^2) \frac{\partial P_l}{\partial x} \quad (4.25)$$

and similarly for the latter half of the equation. we can then write,

$$\begin{aligned} \text{LHS} = & \underbrace{\left[(1-x^2) \frac{dP_l}{dx} P_m \right]_{-1}^1}_0 - \int_{-1}^1 (1-x^2) \frac{dP_m}{dx} \frac{dP_l}{dx} dx \\ & - \underbrace{\left[(1-x^2) \frac{dP_m}{dx} P_l \right]_{-1}^1}_0 - \int_{-1}^1 (1-x^2) \frac{dP_l}{dx} \frac{dP_m}{dx} dx \\ = & 0 \end{aligned} \quad (4.26)$$

We then have that, for $n \neq m$, the LHS is,

$$[m(m+1) - l(l+1)] P_l(x) P_m(x) = 0 \quad (4.27)$$

Furthermore, we can show,

$$\boxed{\int_{-1}^1 P_l(x) P_m(x) dx = \frac{2}{2l+1} \delta_{lm}}. \quad (4.28)$$

4.3.2 Legendre Polynomial Expansion

We can use the Legendre polynomials to perform a Legendre series expansion,

$$f(x) = \sum_{l=0}^{\infty} c_l P_l(x), \quad -1 \leq x \leq 1 \quad (4.29)$$

where the coefficients are given by,

$$c_l = \frac{1}{2}(2l+1) \int_{-1}^1 \int_{-1}^1 f(x) P_l(x) dx. \quad (4.30)$$

4.4 Bessel Functions

We will analyse the temperature distribution on a circular plate. In order to do this, we must solve the diffusion equation,

$$\nabla^2 T = \frac{1}{\alpha^2} \frac{\partial T}{\partial t}. \quad (4.31)$$

Given we have a circular plate, we wish to use the polar form of the Laplacian. We have,

$$\frac{1}{r} \frac{\partial}{\partial r} \left(r \frac{\partial T}{\partial r} \right) + \frac{1}{r^2} \frac{\partial^2 T}{\partial \theta^2} = \frac{1}{\alpha^2} \frac{\partial T}{\partial t}. \quad (4.32)$$

The circular plate has an insulating boundary, and has a cyclic boundary condition,

$$\left. \frac{\partial T}{\partial r} \right|_{r=a} = 0 \quad (4.33)$$

$$T(r, \theta, t) = T(r, \theta + 2n\pi, t). \quad (4.34)$$

We then solve the diffusion equation by separation of variables, $T = R(r)\Theta(\theta)\tau(t)$,

$$\frac{1}{Rr} \frac{d}{dr} \left(r \frac{dR}{dr} \right) + \frac{1}{\Theta r^2} \frac{d^2 \Theta}{d\theta^2} = \frac{1}{\alpha^2 \tau} \frac{d\tau}{dt} = -k^2. \quad (4.35)$$

The *time equation* is,

$$\frac{d\tau}{\tau} = -k^2 \alpha^2 dt \quad (4.36)$$

which has an exponential solution,

$$\tau = A e^{-k^2 \alpha^2 t}. \quad (4.37)$$

Let us analyse the θ and r dependence in eq. (4.35),

$$\frac{r}{R} \left(r \frac{dR}{dr} \right) + k^2 r^2 = -\frac{1}{\Theta} \frac{d^2 \Theta}{d\theta^2} = m^2 \quad (4.38)$$

The cyclic equation is given by,

$$\frac{d\Theta}{d\theta} = -m^2 \Theta \quad (4.39)$$

which has a trigonometric solution,

$$\Theta = A \cos(m\theta) + B \sin(m\theta) \quad (4.40)$$

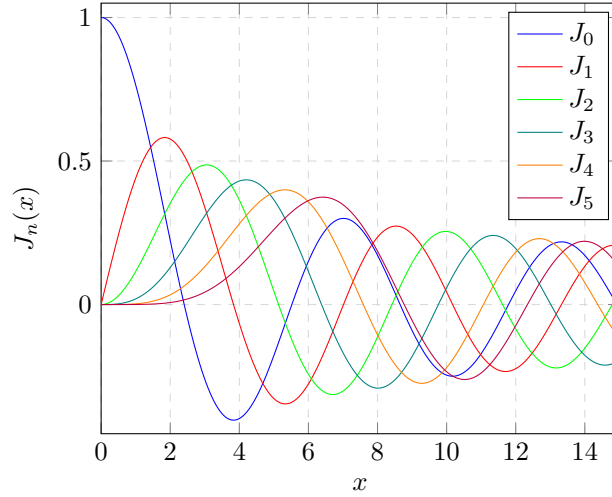


Figure 4.1: The first six Bessel functions.

where we require $m \in \mathbb{N}^+$ to satisfy the cyclic boundary condition.

The *radial equation* is given by,

$$\boxed{r^2 \frac{d^2 R}{dr^2} + r \frac{dR}{dr} + (k^2 r^2 - m^2) R = 0} \quad (4.41)$$

which is also known as *Bessel's equation*. The solutions to this equation is given by,

$$R = J_m(kr) \quad (4.42)$$

where $J_m(kr)$ is an m^{th} order Bessel function.

4.4.1 Solving Bessel Functions

The most general bessel equation is given by,

$$\boxed{x^2 \frac{d^2 y}{dx^2} + x \frac{dy}{dx} + (x^2 - m^2) y = 0} \quad (4.43)$$

We can find a general solution by,

$$y(x) = x^s \sum_{n=0}^{\infty} a_n x^n \quad (4.44)$$

where we impose a boundary condition that $y(0)$ must be finite. Substituting eq (4.44) into eq. (4.43),

$$x \frac{dx}{dy} = \sum_{n=0}^{\infty} a_n (n+s) x^{n+s} \quad (4.45)$$

$$x^2 \frac{d^2 x}{dy^2} = \sum_{n=0}^{\infty} a_n (n+s)(n+s-1) x^{n+s} \quad (4.46)$$

$$(4.47)$$

$$\begin{aligned} (x^2 - m^2)y &= \sum_{n=0}^{\infty} a_n x^{n+s+2} - m^2 \sum_{n=0}^{\infty} a_n x^{n+s} \\ &= \sum_{n=2}^{\infty} a_{n-2} x^{n+s} - m^2 \sum_{n=0}^{\infty} a_n x^{n+s} \end{aligned} \quad (4.48)$$

Putting all these terms together, we can get rid of the x^{n+s} as we require the equation to be true $\forall x$.

$$\sum_{n=0}^{\infty} a_n [(n+s)(n+s-1) + (n+s) - m^2] + \sum_{n=2}^{\infty} a_{n-2} = 0. \quad (4.49)$$

We must first consider $n = 0$ and $n = 1$ before we can find the recurrence relation. For $n = 0$,

$$\begin{aligned} a_0 s(s-1) + a_0 s - m^2 a_0 &= 0 \\ a_0 s^2 - m^2 a_0 & \\ \implies s &= \pm m. \end{aligned} \quad (4.50)$$

For $n = 1$,

$$\begin{aligned} a_1 s(s+1) + a_1(s+1) - m^2 a_1 &= 0 \\ a_1(s+1)(s+1) - m^2 a_1 &= 0 \\ \implies a_1 [(s+1)^2 - m^2] & \end{aligned} \quad (4.51)$$

for which we require $a_1 = 0$ unless $s = \pm m = -\frac{1}{2}$. Otherwise, there will be no odd terms in our solution.

We may now move onto the general recurrence relation,

$$\begin{aligned} [(n+s)(n+s-1) + (n+s) - m^2] a_n - a_{n-2} &= 0 \\ [(n+s)^2 - m^2] a_n + a_{n-2} &= 0 \end{aligned} \quad (4.52)$$

we then have the recurrence relation,

$$a_n = -\frac{a_{n-2}}{(n+m)^2 - m^2} = -\frac{a_{n-2}}{(2m+n)n} \quad (4.53)$$

which can be further generalised to,

$$a_{2j} = (-1)^j \frac{m!}{2^{2j} j! (m+j)!} a_0 \quad (4.54)$$

where $j \in \mathbb{Z}^+$. We can then write the general solution to Bessel's equation,

$$y(x) = \sum_{n=0}^{\infty} a_n x^{m+n} = \sum_{j=0}^{\infty} a_{2j} x^{m+2j}. \quad (4.55)$$

We often rewrite $y(x)$ as,

$$y(x) = a_0 m! 2^m \sum_{j=0}^{\infty} \frac{(-1)^j}{2^{2j} j! (m+j)!} x^{m+2j} \quad (4.56)$$

where we obtain *Bessel's function of the first kind*,

$$\boxed{J_m(x) = \sum_{j=0}^{\infty} \frac{(-1)^j}{j! (m+j)!} \left(\frac{x}{2}\right)^{m+2j}} \quad (4.57)$$

which obeys the orthogonality condition,

$$\int_0^L x J_p\left(\frac{\chi_n x}{l}\right) J_p\left(\frac{\chi_m x}{l}\right) dx = \frac{l^2}{2} [J_{p+1}(\chi_m)]^2 \delta_{mn} \quad (4.58)$$

where,

$$J_p(\chi_n) = 0 \quad n \in \mathbb{Z}^+ \quad (4.59)$$

i.e., χ_n is the n th zero point of the p th Bessel function.

For non-integer values of m , we must use the *gamma function*,

$$J_m(x) = \sum_{j=0}^{\infty} \frac{(-1)^j}{j! \Gamma(m+j+1)} \left(\frac{x}{2}\right)^{m+2j} \quad (4.60)$$

where Γ is defined,

$$\Gamma(s) = \int_0^\infty e^{-t} t^{s-1} dt. \quad (4.61)$$

We find that,

$$J_m(-x) = (-1)^m J_m(x) \quad (4.62)$$

so the m th Bessel function is even for even m , and odd for odd m .

Convergence of the solution

By considering,

$$|r_n| = \left| \frac{a_{n+2}x^{n+2}}{a_nx^n} \right| = \left| \frac{x^2}{(n+2+m)} \right| \quad (4.63)$$

which clearly converges, as in the limit $n \rightarrow \infty$, $|r_n| \rightarrow 0 \forall x$

4.4.2 Modes in a Circular Membrane

Consider a circular membrane clamped at its edges. Using the typical 3D wave equation, we use separation of variables,

$$\phi = F(r, \theta)T(t) \quad (4.64)$$

where $F(r, \theta) = R(r)\Theta(\theta)$. We have,

$$\frac{1}{F} \nabla^2 F = \frac{1}{c^2 T} \frac{\partial^2 T}{\partial t^2} = -k^2 \quad (4.65)$$

and applying the laplacian in polar coordinates,

$$\frac{1}{r} \frac{\partial}{\partial r} \left(r \frac{\partial F}{\partial r} \right) + \frac{1}{r^2} \frac{\partial^2 F}{\partial \theta^2} + k^2 F = 0. \quad (4.66)$$

Let us multiply eq. (4.66) by r^2/F ,

$$\underbrace{\frac{r}{R} \frac{d}{dr} \left(r \frac{dR}{dr} \right)}_{n^2} k^2 r^2 + \underbrace{\frac{1}{\Theta} \frac{d\Theta}{d\theta}}_{-n^2} = 0. \quad (4.67)$$

We then have,

$$\frac{1}{\Theta} \frac{d\Theta}{d\theta} + n^2 = 0 \quad (4.68)$$

$$\frac{r}{R} \frac{d}{dr} \left(r \frac{dR}{dr} \right) + k^2 r^2 - n^2 = 0 \quad (4.69)$$

We can clearly see that eq. (4.68) produces a cyclic solution. However, rearranging eq. (4.69),

$$r \frac{d}{dr} \left(r \frac{dR}{dr} \right) + (k^2 r^2 - n^2) R = 0 \quad (4.70)$$

which is Bessel's equation, so, $R(r) \approx J_n(kr)$. Given the boundary condition,

$$R(a) = 0 \quad (4.71)$$

we have,

$$J_n(ka) = 0 \quad (4.72)$$

so we require,

$$k_{nm}a = \chi_{nm} \quad (4.73)$$

where χ_{nm} is the m th root of the n th Bessel function of the first kind. We then require the angular frequency of the modes of oscillation of the membrane to be,

$$\omega_{nm} = k_{nm}c = \frac{\chi_{nm}c}{a}. \quad (4.74)$$

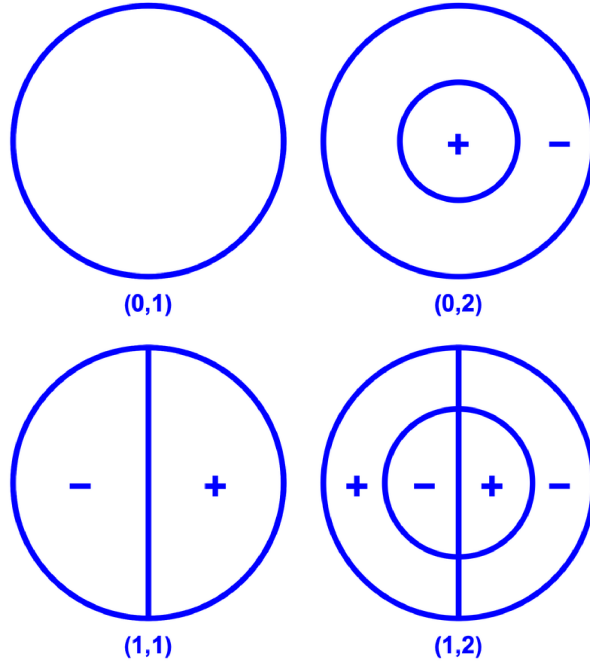


Figure 4.2

Some normal modes of the membrane are shown in figure 4.2.

NOTE: Frequencies of vibration are not integer multiples of the fundamental mode, as it might be with a regular wave.

The final solution for the vibrations of the membrane is given by,

$$\phi = \sum_{n=0}^{\infty} \sum_{m=0}^{\infty} J_n(k_{nm}r) [A_{nm} \cos(n\theta) + B_{nm} \sin(n\theta)] \cos(\omega_{nm}t) \quad (4.75)$$

4.5 Spherical Waves

The spherical wave equation is governed by the laplacian in spherical polar coordinates. The spherical wave equation is given by,

$$\frac{1}{r^2} \frac{\partial}{\partial r} \left(r^2 \frac{\partial f}{\partial r} \right) + \frac{1}{r^2 \sin \theta} \frac{\partial}{\partial \theta} \left(\sin \theta \frac{\partial f}{\partial \theta} \right) + \frac{1}{r^2 \sin^2 \theta} \frac{\partial^2 f}{\partial \phi^2} = \frac{1}{c^2} \frac{\partial^2 f}{\partial t^2}. \quad (4.76)$$

Let us for now consider the solution on the surface, $r = a$. Let us use a separation of variables,

$$F(\theta, \phi, t) = Y(\theta, \phi)T(t) \quad (4.77)$$

We have,

$$\frac{1}{Y a^2 \sin \theta} \frac{\partial}{\partial \theta} \left(\sin \theta \frac{\partial Y}{\partial \theta} \right) + \frac{1}{Y a^2 \sin^2 \theta} \frac{\partial^2 Y}{\partial \phi^2} = \frac{1}{c^2 T} \frac{d^2 T}{dt^2} = -\frac{l(l+1)}{a^2}. \quad (4.78)$$

Analysing the time equation,

$$\frac{d^2 T}{dt^2} = -\frac{l(l+1)}{a^2 c^2} T \quad (4.79)$$

we have a trivial SHM solution,

$$T = A \sin(\omega t) + B \cos(\omega t) \quad (4.80)$$

where

$$\omega^2 = \frac{c^2}{a^2} l(l+1). \quad (4.81)$$

For the angular solution, we use separation of variables,

$$Y = \Theta(\theta)\Phi(\phi) \quad (4.82)$$

to obtain,

$$\begin{aligned} \frac{1}{\sin\theta} \frac{d}{d\theta} \left(\sin\theta \frac{d\Theta}{d\theta} \right) + \frac{1}{\sin^2\theta} \frac{d^2\Phi}{d\phi^2} = -l(l+1)\Theta\Phi \\ \underbrace{\frac{\sin\theta}{\Theta} \frac{d}{d\theta} (\sin\theta) + l(l+1)\sin^2\theta}_{m^2} + \underbrace{\frac{1}{\Phi} \frac{d^2\Phi}{d\phi^2}}_{-m^2} = 0 \end{aligned} \quad (4.83)$$

Let us consider a solution of constant ϕ , and set $m = 0$. Our equation is then purely in θ ,

$$\frac{\sin\theta}{\Theta} \frac{d}{d\theta} \left(\sin\theta \frac{d\Theta}{d\theta} \right) + l(l+1)\sin^2\theta = 0. \quad (4.84)$$

We can solve this by a substitution $x = \cos\theta$,

$$\frac{dx}{d\theta} = -\sin\theta \quad \implies \quad d\theta = -\frac{dx}{\sin\theta} \quad (4.85)$$

Our equation then becomes,

$$-\frac{\sin^2\theta}{\Theta} \frac{d}{dx} \left(-\sin^2\theta \frac{d\Theta}{dx} \right) + l(l+1)\sin^2\theta = 0 \quad (4.86)$$

If we note that $\sin^2\theta = 1 - x^2$,

$$\frac{d}{dx} \left((1 - x^2) \frac{d\Theta}{dx} \right) + l(l+1)\Theta = 0 \quad (4.87)$$

which is Legendre's polynomials. Our solution is then,

$$\Theta = P_l(x). \quad (4.88)$$

The general solution, considering only θ and t dependence is then,

$$f(\theta, t) = \sum_{l=0}^{\infty} P_l(\cos\theta) (A_l \cos\omega t + B_l \sin\omega t) \quad (4.89)$$

4.5.1 ϕ Dependent Solution

Let us now consider solutions where $m \neq 0$. Our angular equations become,

$$\frac{1}{\Phi} \frac{d^2\Phi}{d\phi^2} = -m^2 \quad (4.90)$$

$$\frac{1}{\sin\theta} \frac{d}{d\theta} \left(\sin^2\theta \frac{d\Theta}{d\theta} \right) + \left[l(l+1) - \frac{m^2}{\sin^2\theta} \right] \Theta = 0 \quad (4.91)$$

Eq. (4.90) is trivial,

$$\Phi = A \sin m\phi + B \cos m\phi, \quad m \in \mathbb{Z}^+. \quad (4.92)$$

Eq. (4.91) can be solved if we consider,

$$x = \cos\theta \quad (4.93)$$

$$\frac{dx}{d\theta} = -\sin\theta \therefore d\theta = \frac{dx}{\sin\theta} \quad (4.94)$$

we can rewrite eq. (4.91) as,

$$\frac{d}{dx} \left((1 - x^2) \frac{d\Theta}{dx} \right) + \left[l(l+1) - \frac{m^2}{1 - x^2} \right] \Theta = 0. \quad (4.95)$$

For $m = 0$, we obtain Legendre polynomial solutions. However, for $m \neq 0$, we obtain associated Legendre polynomials,

$$\Theta = P_l^m(x), \quad |m| \leq l. \quad (4.96)$$

We then find the total angular solution is,

$$Y(\theta, \phi) = P_l^m(\cos\theta) [A \sin m\phi + B \cos m\phi] \quad (4.97)$$

4.5.2 Waves Inside the Sphere (r dependence)

If we take into account r dependence, we retain the angular solutions from before. Focusing on the radial and time solutions,

$$\frac{1}{Rr^2} \frac{d}{dr} \left(r^2 \frac{dR}{dr} \right) - \frac{l(l+1)}{r^2} = \frac{1}{c^2 T} \frac{d^2 T}{dt^2} = -k^2. \quad (4.98)$$

We see the time solutions are SHM. Focussing on the radial equation,

$$\frac{d}{dr} \left(r^2 \frac{dR}{dr} \right) + [k^2 r^2 - l(l+1)] R = 0. \quad (4.99)$$

Let us recall the boundary condition $R(a) = 0$, where a is the surface of the sphere. Let $x = kr$ to rewrite eq. (4.98),

$$\boxed{x^2 \frac{d^2 R}{dx^2} + 2x \frac{dR}{dx} + [x^2 - l(l+1)] R = 0} \quad (4.100)$$

which is the form of the *spherical Bessel equation*. The solution to the spherical Bessel equation are the spherical bessel functions,

$$R = j_l(kr). \quad (4.101)$$

It can be shown that,

$$j_l(x) = \sqrt{\frac{\pi}{2x}} J_{l+\frac{1}{2}}(x). \quad (4.102)$$

Applying the boundary condition,

$$j_l(k_{lm}a) = 0. \quad (4.103)$$

The final solution to the spherical wave equation is given by,

$$f = \sum_{l,m,n}^{\infty} j_l(k_{ln}r) P_l^m(\cos \theta) [\sin(m\phi) + \cos(m\phi)] [A_{lmn} \cos(\omega_{lmn}t) + B_{lmn} \sin(\omega_{lmn}t)] \quad (4.104)$$

Appendix A

Examples: Differential Equations

A.1 1D Wave Equation

The 1 dimensional wave equation for a wavefunction ϕ is given by,

$$\boxed{\frac{\partial^2 \phi}{\partial x^2} = \frac{1}{c^2} \frac{\partial^2 \phi}{\partial t^2}}. \quad (\text{A.1})$$

A.1.1 Euler/d'Alembert solution

We can find a general solution to eq. (A.1) by using the substitution,

$$v = x - ct \Leftarrow \text{Backward component} \quad (\text{A.2})$$

$$u = x + ct \Leftarrow \text{Forward component}. \quad (\text{A.3})$$

Computing the derivative with respect to x by the chain rule,

$$\frac{\partial \phi}{\partial x} = \frac{\partial \phi}{\partial v} \frac{\partial v}{\partial x} + \frac{\partial \phi}{\partial u} \frac{\partial u}{\partial x} = \left(\frac{\partial}{\partial v} + \frac{\partial}{\partial u} \right) \phi \quad (\text{A.4})$$

from which the second derivative follows,

$$\frac{\partial^2 \phi}{\partial x^2} = \left(\frac{\partial}{\partial u} + \frac{\partial}{\partial v} \right)^2 \phi. \quad (\text{A.5})$$

Similarly for the time component,

$$\frac{\partial^2 \phi}{\partial t^2} = c^2 \left(\frac{\partial}{\partial u} - \frac{\partial}{\partial v} \right)^2 \phi. \quad (\text{A.6})$$

Applying eqs. (A.5) and (A.6) to eq. (A.1), we find,

$$\begin{aligned} \left(\frac{\partial^2}{\partial u^2} + \frac{\partial^2}{\partial v^2} + 2 \frac{\partial}{\partial u} \frac{\partial}{\partial v} \right) &= \frac{c^2}{c^2} \left(\frac{\partial^2}{\partial u^2} + \frac{\partial^2}{\partial v^2} - 2 \frac{\partial}{\partial u} \frac{\partial}{\partial v} \right) \\ \Rightarrow \left(\frac{\partial}{\partial u} \frac{\partial}{\partial v} \right) \phi &= 0 \Rightarrow \text{The solution is a sum of backward and forward components.} \end{aligned} \quad (\text{A.7})$$

Thus the general solution to eq. (A.1) is,

$$\boxed{\phi = \phi(x + ct) + \phi(x - ct)}. \quad (\text{A.8})$$

A.2 Laplace's Equation

Laplace's equation is given by,

$$\boxed{\nabla^2 \phi = 0} \quad (\text{A.9})$$

and can be readily solved using separation of variables.

A.3 Diffusion Equation

The diffusion equation is given by,

$$\boxed{\nabla^2 T = \frac{1}{\alpha} \frac{\partial T}{\partial t}} \quad (\text{A.10})$$

which reduces to Laplace's equation for a steady-state system.

Appendix B

Misc. Notes

B.1 Plane Wave Nature

We have that travelling waves in 3D are given by $e^{i(\mathbf{k} \cdot \mathbf{x})}$ and have a plane wave nature, such that,

$$(\mathbf{k} \cdot \mathbf{x} - \omega t) = \phi \tag{B.1}$$

where ϕ is constant. This corresponds to a plane wave with constant phase. So, we can write,

$$\begin{aligned} |\mathbf{k}||\mathbf{x}| - \omega t &= \phi \\ \implies |\mathbf{x}| &= \frac{\omega t + \phi}{|\mathbf{k}|} \end{aligned} \tag{B.2}$$

which is the equation of a sphere with radius $\frac{\omega t + \phi}{|\mathbf{k}|}$. Thus, a travelling wave in 3D represents spherical waves propagating outwards with velocity $\frac{\omega}{|\mathbf{k}|}$.