

electro magnetic field theory

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ELECTROMAGNETIC FIELD THEORY

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EXERCISES**

by

Tobia Carozzi, Anders Eriksson, Bengt Lundborg,
Bo Thidé and Mattias Waldenvik

ELECTROMAGNETIC FIELD THEORY

Bo Thidé

*Department of Space and Plasma Physics
Uppsala University*

and

*Swedish Institute of Space Physics
Uppsala Division
Sweden*



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To the memory of
LEV MIKHAILOVICH ERUKHIMOV
dear friend, remarkable physicist
and a truly great human.

This book is the result of a twenty-five year long love affair. In 1972, I took my first advanced course in electrodynamics at the Theoretical Physics department, Uppsala University. Shortly thereafter, I joined the research group there and took on the task of helping my supervisor, professor PER-OLOF FRÖMAN, with the preparation of a new version of his lecture notes on Electricity Theory. These two things opened up my eyes for the beauty and intricacy of electrodynamics, already at the classical level, and I fell in love with it.

Ever since that time, I have off and on had reason to return to electrodynamics, both in my studies, research and teaching, and the current book is the result of my own teaching of a course in advanced electrodynamics at Uppsala University some twenty odd years after I experienced the first encounter with this subject. The book is the outgrowth of the lecture notes that I prepared for the four-credit course Electrodynamics that was introduced in the Uppsala University curriculum in 1992, to become the five-credit course Classical Electrodynamics in 1997. To some extent, parts of these notes were based on lecture notes prepared, in Swedish, by BENGT LUNDBORG who created, developed and taught the earlier, two-credit course Electromagnetic Radiation at our faculty.

Intended primarily as a textbook for physics students at the advanced undergraduate or beginning graduate level, I hope the book may be useful for research workers too. It provides a thorough treatment of the theory of electrodynamics, mainly from a classical field theoretical point of view, and includes such things as electrostatics and magnetostatics and their unification into electrodynamics, the electromagnetic potentials, gauge transformations, covariant formulation of classical electrodynamics, force, momentum and energy of the electromagnetic field, radiation and scattering phenomena, electromagnetic waves and their propagation in vacuum and in media, and covariant Lagrangian/Hamiltonian field theoretical methods for electromagnetic fields, particles and interactions. The aim has been to write a book that can serve both as an advanced text in Classical Electrodynamics and as a preparation for studies in Quantum Electrodynamics and related subjects.

In an attempt to encourage participation by other scientists and students in the authoring of this book, and to ensure its quality and scope to make it useful in higher university education anywhere in the world, it was produced within a World-Wide Web (WWW) project. This turned out to be a rather successful move. By making an electronic version of the book freely down-loadable on the net, I have not only received comments on it from fellow Internet physicists around the world, but know, from WWW ‘hit’ statistics that at the time of writing this, the book serves as a frequently used Internet resource. This way it is my hope that it will be particularly useful for students and researchers working under financial or other circumstances that make it difficult to procure a printed copy of the book.

I am grateful not only to Per-Olof Fröman and Bengt Lundborg for providing the inspiration for my writing this book, but also to CHRISTER WAHLBERG at Uppsala University for interesting discussions on electrodynamics in general and on this book in particular, and to my former graduate students MATTIAS WALDENVIK and TOBIA CAROZZI as well as ANDERS ERIKSSON, all at the Swedish Institute of Space Physics, Uppsala Division, and who have participated in the teaching and commented on the material covered in the course and in this book. Thanks are also due to my long-term space physics colleague HELMUT KOPKA of the Max-Planck-Institut für Aeronomie, Lindau, Germany, who not only taught me about the practical aspects of the of high-power radio wave transmitters and transmission lines, but also about the more delicate aspects of typesetting a book in $\text{T}_{\text{E}}\text{X}$ and $\text{L}_{\text{A}}\text{T}_{\text{E}}\text{X}$. I am particularly indebted to Academician professor VITALIY L. GINZBURG for his many fascinating and very elucidating lectures, comments and historical footnotes on electromagnetic radiation while cruising on the Volga river during our joint Russian-Swedish summer schools.

Finally, I would like to thank all students and Internet users who have downloaded and commented on the book during its life on the World-Wide Web.

I dedicate this book to my son MATTIAS, my daughter KAROLINA, my high-school physics teacher, STAFFAN RÖSBY, and to my fellow members of the CAPELLA PEDAGOGICA UPSALIENSIS.

Uppsala, Sweden
November, 1999

BO THIDÉ

Classical Electrodynamics

Classical electrodynamics deals with electric and magnetic fields and interactions caused by *macroscopic* distributions of electric charges and currents. This means that the concepts of localised charges and currents assume the validity of certain mathematical limiting processes in which it is considered possible for the charge and current distributions to be localised in infinitesimally small volumes of space. This is in obvious contradiction to electromagnetism on a *microscopic* scale, where charges and currents are known to be spatially extended objects. However, the limiting processes yield results which are correct on a macroscopic scale.

In this Chapter we start with the force interactions in classical electrostatics and classical magnetostatics and introduce the static electric and magnetic fields and find two uncoupled systems of equations for them. Then we see how the conservation of electric charge and its relation to electric current leads to the dynamic connection between electricity and magnetism and how the two can be unified in one theory, classical electrodynamics, described by one system of coupled dynamic field equations.

1.1 Electrostatics

The theory that describes physical phenomena related to the interaction between stationary electric charges or charge distributions in space is called *electrostatics*.

1.1.1 Coulomb's law

It has been found experimentally that in classical electrostatics the interaction between two stationary electrically charged bodies can be described in terms of a

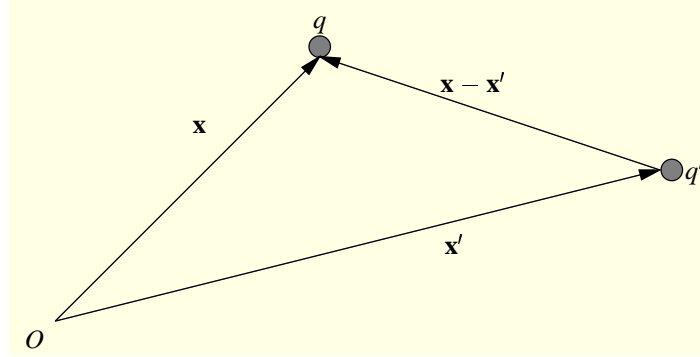


Figure 1.1. Coulomb's law describes how a static electric charge q , located at a point \mathbf{x} relative to the origin O , experiences an electrostatic force from a static electric charge q' located at \mathbf{x}' .

mechanical force. Let us consider the simple case depicted in figure 1.1 where \mathbf{F} denotes the force acting on a charged particle with charge q located at \mathbf{x} , due to the presence of a charge q' located at \mathbf{x}' . According to *Coulomb's law* this force is, in vacuum, given by the expression

$$\begin{aligned}\mathbf{F}(\mathbf{x}) &= \frac{qq'}{4\pi\epsilon_0} \frac{\mathbf{x} - \mathbf{x}'}{|\mathbf{x} - \mathbf{x}'|^3} \\ &= -\frac{qq'}{4\pi\epsilon_0} \nabla \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right)\end{aligned}\quad (1.1)$$

where we have used results from Example M.6 on page 186. In SI units, which we shall use throughout, the force \mathbf{F} is measured in Newton (N), the charges q and q' in Coulomb (C) [= Ampère-seconds (As)], and the length $|\mathbf{x} - \mathbf{x}'|$ in metres (m). The constant $\epsilon_0 = 10^7/(4\pi c^2) \approx 8.8542 \times 10^{-12}$ Farad per metre (F/m) is the *vacuum permittivity* and $c \approx 2.9979 \times 10^8$ m/s is the speed of light in vacuum. In CGS units $\epsilon_0 = 1/(4\pi)$ and the force is measured in dyne, the charge in statcoulomb, and length in centimetres (cm).

1.1.2 The electrostatic field

Instead of describing the electrostatic interaction in terms of a “force action at a distance,” it turns out that it is often more convenient to introduce the concept of a field and to describe the electrostatic interaction in terms of a static vectorial

electric field \mathbf{E}^{stat} defined by the limiting process

$$\mathbf{E}^{\text{stat}} \stackrel{\text{def}}{=} \lim_{q \rightarrow 0} \frac{\mathbf{F}}{q} \quad (1.2)$$

where \mathbf{F} is the electrostatic force, as defined in equation (1.1) on the facing page, from a net charge q' on the test particle with a small electric net charge q . Since the purpose of the limiting process is to assure that the test charge q does not influence the field, the expression for \mathbf{E}^{stat} does not depend explicitly on q but only on the charge q' and the relative radius vector $\mathbf{x} - \mathbf{x}'$. This means that we can say that any net electric charge produces an electric field in the space that surrounds it, regardless of the existence of a second charge anywhere in this space.¹

Using formulae (1.1) and (1.2), we find that the electrostatic field \mathbf{E}^{stat} at the field point \mathbf{x} (also known as the *observation point*), due to a field-producing charge q' at the source point \mathbf{x}' , is given by

$$\begin{aligned} \mathbf{E}^{\text{stat}}(\mathbf{x}) &= \frac{q'}{4\pi\epsilon_0} \frac{\mathbf{x} - \mathbf{x}'}{|\mathbf{x} - \mathbf{x}'|^3} \\ &= -\frac{q'}{4\pi\epsilon_0} \nabla \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right) \end{aligned} \quad (1.3)$$

In the presence of several field producing discrete charges q'_i , at \mathbf{x}'_i , $i = 1, 2, 3, \dots$, respectively, the assumption of linearity of vacuum² allows us to superimpose their individual \mathbf{E} fields into a total \mathbf{E} field

$$\mathbf{E}^{\text{stat}}(\mathbf{x}) = \sum_i \frac{q'_i}{4\pi\epsilon_0} \frac{\mathbf{x} - \mathbf{x}'_i}{|\mathbf{x} - \mathbf{x}'_i|^3} \quad (1.4)$$

If the discrete charges are small and numerous enough, we introduce the *charge*

¹In the preface to the first edition of the first volume of his book *A Treatise on Electricity and Magnetism*, first published in 1873, James Clerk Maxwell describes this in the following, almost poetic, manner [?]:

“For instance, Faraday, in his mind’s eye, saw lines of force traversing all space where the mathematicians saw centres of force attracting at a distance: Faraday saw a medium where they saw nothing but distance: Faraday sought the seat of the phenomena in real actions going on in the medium, they were satisfied that they had found it in a power of action at a distance impressed on the electric fluids.”

²In fact, vacuum exhibits a *quantum mechanical nonlinearity* due to *vacuum polarisation effects* manifesting themselves in the momentary creation and annihilation of electron-positron pairs, but classically this nonlinearity is negligible.

density ρ located at \mathbf{x}' and write the total field as

$$\begin{aligned}\mathbf{E}^{\text{stat}}(\mathbf{x}) &= \frac{1}{4\pi\epsilon_0} \int_V \rho(\mathbf{x}') \frac{\mathbf{x} - \mathbf{x}'}{|\mathbf{x} - \mathbf{x}'|^3} d^3x' \\ &= -\frac{1}{4\pi\epsilon_0} \int_V \rho(\mathbf{x}') \nabla \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right) d^3x'\end{aligned}\quad (1.5)$$

where, in the last step, we used formula equation (M.65) on page 186. We emphasize that equation (1.5) above is valid for an arbitrary distribution of charges, including discrete charges, in which case ρ can be expressed in terms of one or more Dirac delta functions.

Since, according to formula equation (M.75) on page 188, $\nabla \times [\nabla\alpha(\mathbf{x})] \equiv \mathbf{0}$ for any 3D \mathbb{R}^3 scalar field $\alpha(\mathbf{x})$, we immediately find that in electrostatics

$$\begin{aligned}\nabla \times \mathbf{E}^{\text{stat}}(\mathbf{x}) &= -\frac{1}{4\pi\epsilon_0} \nabla \times \int_V \rho(\mathbf{x}') \left[\nabla \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right) \right] d^3x' \\ &= -\frac{1}{4\pi\epsilon_0} \int_V \rho(\mathbf{x}') \nabla \times \left[\nabla \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right) \right] d^3x' \\ &= \mathbf{0}\end{aligned}\quad (1.6)$$

I.e., \mathbf{E}^{stat} is an *irrotational* field.

Taking the divergence of the general \mathbf{E}^{stat} expression for an arbitrary charge distribution, equation (1.5), and using the representation of the Dirac delta function, equation (M.70) on page 187, we find that

$$\begin{aligned}\nabla \cdot \mathbf{E}^{\text{stat}}(\mathbf{x}) &= \nabla \cdot \frac{1}{4\pi\epsilon_0} \int_V \rho(\mathbf{x}') \frac{\mathbf{x} - \mathbf{x}'}{|\mathbf{x} - \mathbf{x}'|^3} d^3x' \\ &= -\frac{1}{4\pi\epsilon_0} \int_V \rho(\mathbf{x}') \nabla \cdot \nabla \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right) d^3x' \\ &= -\frac{1}{4\pi\epsilon_0} \int_V \rho(\mathbf{x}') \nabla^2 \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right) d^3x' \\ &= \frac{1}{\epsilon_0} \int_V \rho(\mathbf{x}') \delta(\mathbf{x} - \mathbf{x}') d^3x' \\ &= \frac{\rho(\mathbf{x})}{\epsilon_0}\end{aligned}\quad (1.7)$$

which is *Gauss's law* in differential form.

1.2 Magnetostatics

While electrostatics deals with static charges, *magnetostatics* deals with stationary currents, *i.e.*, charges moving with constant speeds, and the interaction between these currents.

1.2.1 Ampère's law

Experiments on the interaction between two small current loops have shown that they interact via a mechanical force, much the same way that charges interact. Let \mathbf{F} denote such a force acting on a small loop C carrying a current J located at \mathbf{x} , due to the presence of a small loop C' carrying a current J' located at \mathbf{x}' . According to *Ampère's law* this force is, in vacuum, given by the expression

$$\begin{aligned}\mathbf{F}(\mathbf{x}) &= \frac{\mu_0 J J'}{4\pi} \oint_C \oint_{C'} d\mathbf{l} \times \frac{d\mathbf{l}' \times (\mathbf{x} - \mathbf{x}')}{|\mathbf{x} - \mathbf{x}'|^3} \\ &= -\frac{\mu_0 J J'}{4\pi} \oint_C \oint_{C'} d\mathbf{l} \times \left[d\mathbf{l}' \times \nabla \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right) \right]\end{aligned}\tag{1.8}$$

Here $d\mathbf{l}$ and $d\mathbf{l}'$ are tangential line elements of the loops C and C' , respectively, and, in SI units, $\mu_0 = 4\pi \times 10^{-7} \approx 1.2566 \times 10^{-6}$ H/m is the *vacuum permeability*. From the definition of ϵ_0 and μ_0 (in SI units) we observe that

$$\epsilon_0 \mu_0 = \frac{10^7}{4\pi c^2} (\text{F/m}) \times 4\pi \times 10^{-7} (\text{H/m}) = \frac{1}{c^2} (\text{s}^2/\text{m}^2)\tag{1.9}$$

which is a useful relation.

At first glance, equation (1.8) above appears to be unsymmetric in terms of the loops and therefore to be a force law which is in contradiction with Newton's third law. However, by applying the vector triple product "bac-cab" formula (F.56) on page 167, we can rewrite (1.8) in the following way

$$\begin{aligned}\mathbf{F}(\mathbf{x}) &= -\frac{\mu_0 J J'}{4\pi} \oint_C \oint_{C'} \left[d\mathbf{l} \cdot \nabla \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right) \right] d\mathbf{l}' \\ &\quad - \frac{\mu_0 J J'}{4\pi} \oint_C \oint_{C'} \frac{\mathbf{x} - \mathbf{x}'}{|\mathbf{x} - \mathbf{x}'|^3} d\mathbf{l} \cdot d\mathbf{l}'\end{aligned}\tag{1.10}$$

Recognising the fact the integrand in the first integral is an exact differential so that this integral vanishes, we can rewrite the force expression, equation (1.8) above, in

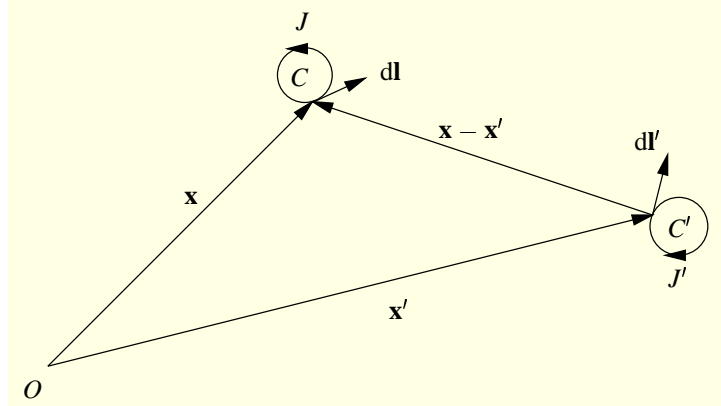


Figure 1.2. Ampère's law describes how a small loop C , carrying a static electric current J through its tangential line element $d\mathbf{l}$ located at \mathbf{x} , experiences a magnetostatic force from a small loop C' , carrying a static electric current J' through the tangential line element $d\mathbf{l}'$ located at \mathbf{x}' . The loops can have arbitrary shapes as long as they are simple and closed.

the following symmetric way

$$\mathbf{F}(\mathbf{x}) = -\frac{\mu_0 J J'}{4\pi} \oint_C \oint_{C'} \frac{\mathbf{x} - \mathbf{x}'}{|\mathbf{x} - \mathbf{x}'|^3} d\mathbf{l} \cdot d\mathbf{l}' \quad (1.11)$$

This clearly exhibits the expected symmetry in terms of loops C and C' .

1.2.2 The magnetostatic field

In analogy with the electrostatic case, we may attribute the magnetostatic interaction to a vectorial *magnetic field* \mathbf{B}^{stat} . It turns out that \mathbf{B}^{stat} can be defined through

$$d\mathbf{B}^{\text{stat}}(\mathbf{x}) \stackrel{\text{def}}{=} \frac{\mu_0 J'}{4\pi} d\mathbf{l}' \times \frac{\mathbf{x} - \mathbf{x}'}{|\mathbf{x} - \mathbf{x}'|^3} \quad (1.12)$$

which expresses the small element $d\mathbf{B}^{\text{stat}}(\mathbf{x})$ of the static magnetic field set up at the field point \mathbf{x} by a small line element $d\mathbf{l}'$ of stationary current J' at the source point \mathbf{x}' . The SI unit for the magnetic field, sometimes called the *magnetic flux density* or *magnetic induction*, is Tesla (T).

If we generalise expression (1.12) to an integrated steady state current distribu-

tion $\mathbf{j}(\mathbf{x})$, we obtain *Biot-Savart's law*:

$$\begin{aligned}\mathbf{B}^{\text{stat}}(\mathbf{x}) &= \frac{\mu_0}{4\pi} \int_V \mathbf{j}(\mathbf{x}') \times \frac{\mathbf{x} - \mathbf{x}'}{|\mathbf{x} - \mathbf{x}'|^3} d^3x' \\ &= -\frac{\mu_0}{4\pi} \int_V \mathbf{j}(\mathbf{x}') \times \nabla \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right) d^3x'\end{aligned}\quad (1.13)$$

Comparing equation (1.5) on page 4 with equation (1.13), we see that there exists a close analogy between the expressions for \mathbf{E}^{stat} and \mathbf{B}^{stat} but that they differ in their vectorial characteristics. With this definition of \mathbf{B}^{stat} , equation (1.8) on page 5 may be written

$$\mathbf{F}(\mathbf{x}) = J \oint_C d\mathbf{l} \times \mathbf{B}^{\text{stat}}(\mathbf{x}) \quad (1.14)$$

In order to assess the properties of \mathbf{B}^{stat} , we determine its divergence and curl. Taking the divergence of both sides of equation (1.13) and utilising formula (F.63) on page 167, we obtain

$$\begin{aligned}\nabla \cdot \mathbf{B}^{\text{stat}}(\mathbf{x}) &= -\frac{\mu_0}{4\pi} \nabla \cdot \int_V \mathbf{j}(\mathbf{x}') \times \nabla \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right) d^3x' \\ &= -\frac{\mu_0}{4\pi} \int_V \nabla \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right) \cdot [\nabla \times \mathbf{j}(\mathbf{x}')] d^3x' \\ &\quad + \frac{\mu_0}{4\pi} \int_V \mathbf{j}(\mathbf{x}') \cdot \left[\nabla \times \nabla \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right) \right] d^3x' \\ &= 0\end{aligned}\quad (1.15)$$

where the first term vanishes because $\mathbf{j}(\mathbf{x}')$ is independent of \mathbf{x} so that $\nabla \times \mathbf{j}(\mathbf{x}') \equiv \mathbf{0}$, and the second term vanishes since, according to equation (M.75) on page 188, $\nabla \times [\nabla \alpha(\mathbf{x})]$ vanishes for any scalar field $\alpha(\mathbf{x})$.

Applying the operator “bac-cab” rule, formula (F.69) on page 168, the curl of equation (1.13) above can be written

$$\begin{aligned}\nabla \times \mathbf{B}^{\text{stat}}(\mathbf{x}) &= -\frac{\mu_0}{4\pi} \nabla \times \int_V \mathbf{j}(\mathbf{x}') \times \nabla \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right) d^3x' \\ &= -\frac{\mu_0}{4\pi} \int_V \mathbf{j}(\mathbf{x}') \nabla^2 \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right) d^3x' \\ &\quad + \frac{\mu_0}{4\pi} \int_V [\mathbf{j}(\mathbf{x}') \cdot \nabla'] \nabla' \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right) d^3x'\end{aligned}\quad (1.16)$$

If, in the first of the two integrals on the right hand side, we use the representation of the Dirac delta function equation (M.70) on page 187, and integrate the second

one by parts, by utilising formula (F.61) on page 167 as follows:

$$\begin{aligned}
& \int_V [\mathbf{j}(\mathbf{x}') \cdot \nabla'] \nabla' \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right) d^3x' \\
&= \hat{\mathbf{x}}_k \int_V \nabla' \cdot \left\{ \mathbf{j}(\mathbf{x}') \left[\frac{\partial}{\partial x'_k} \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right) \right] \right\} d^3x' \\
&\quad - \int_V [(\nabla' \cdot \mathbf{j}(\mathbf{x}')) \nabla' \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right)] d^3x' \\
&= \hat{\mathbf{x}}_k \int_S \mathbf{j}(\mathbf{x}') \frac{\partial}{\partial x'_k} \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right) \cdot d\mathbf{S} \\
&\quad - \int_V [(\nabla' \cdot \mathbf{j}(\mathbf{x}')) \nabla' \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right)] d^3x' = \mathbf{0}
\end{aligned} \tag{1.17}$$

Here the first integral, obtained by applying Gauss's theorem, vanishes when integrated over a large sphere far away from the localised source $\mathbf{j}(\mathbf{x}')$, and the second one vanishes because $\nabla \cdot \mathbf{j} = 0$ for stationary currents (no charge accumulation in space). The net result is simply

$$\nabla \times \mathbf{B}^{\text{stat}}(\mathbf{x}) = \mu_0 \int_V \mathbf{j}(\mathbf{x}') \delta(\mathbf{x} - \mathbf{x}') d^3x' = \mu_0 \mathbf{j}(\mathbf{x}) \tag{1.18}$$

1.3 Electrodynamics

As we saw in the previous sections, the laws of electrostatics and magnetostatics can be summarised in two pairs of time-independent, uncoupled vector differential equations, namely the *equations of classical electrostatics*

$$\nabla \cdot \mathbf{E}^{\text{stat}}(\mathbf{x}) = \frac{\rho(\mathbf{x})}{\epsilon_0} \tag{1.19a}$$

$$\nabla \times \mathbf{E}^{\text{stat}}(\mathbf{x}) = \mathbf{0} \tag{1.19b}$$

and the *equations of classical magnetostatics*

$$\nabla \cdot \mathbf{B}^{\text{stat}}(\mathbf{x}) = 0 \tag{1.20a}$$

$$\nabla \times \mathbf{B}^{\text{stat}}(\mathbf{x}) = \mu_0 \mathbf{j}(\mathbf{x}) \tag{1.20b}$$

Since there is nothing *a priori* which connects \mathbf{E}^{stat} directly with \mathbf{B}^{stat} , we must consider classical electrostatics and classical magnetostatics as two independent theories.

However, when we include time-dependence, these theories are unified into one theory, *classical electrodynamics*. This unification of the theories of electricity and magnetism is motivated by two empirically established facts:

1. Electric charge is a conserved quantity and current is a transport of electric charge. This fact manifests itself in the equation of continuity and, as a consequence, in Maxwell's displacement current.
2. A change in the magnetic flux through a loop will induce an EMF electric field in the loop. This is the celebrated Faraday's law of induction.

1.3.1 Equation of continuity

Let \mathbf{j} denote the electric current density (A/m^2). In the simplest case it can be defined as $\mathbf{j} = \mathbf{v}\rho$ where \mathbf{v} is the velocity of the charge density. In general, \mathbf{j} has to be defined in statistical mechanical terms as $\mathbf{j}(t, \mathbf{x}) = \sum_{\alpha} q_{\alpha} \int \mathbf{v} f_{\alpha}(t, \mathbf{x}, \mathbf{v}) d^3v$ where $f_{\alpha}(t, \mathbf{x}, \mathbf{v})$ is the (normalised) distribution function for particle species α with electrical charge q_{α} .

The *electric charge conservation law* can be formulated in the *equation of continuity*

$$\frac{\partial \rho(t, \mathbf{x})}{\partial t} + \nabla \cdot \mathbf{j}(t, \mathbf{x}) = 0 \quad (1.21)$$

which states that the time rate of change of electric charge $\rho(t, \mathbf{x})$ is balanced by a divergence in the electric current density $\mathbf{j}(t, \mathbf{x})$.

1.3.2 Maxwell's displacement current

We recall from the derivation of equation (1.18) on the preceding page that there we used the fact that in magnetostatics $\nabla \cdot \mathbf{j}(\mathbf{x}) = 0$. In the case of non-stationary sources and fields, we must, in accordance with the continuity equation (1.21), set $\nabla \cdot \mathbf{j}(t, \mathbf{x}) = -\partial \rho(t, \mathbf{x})/\partial t$. Doing so, and formally repeating the steps in the derivation of equation (1.18) on the preceding page, we would obtain the formal result

$$\begin{aligned} \nabla \times \mathbf{B}(t, \mathbf{x}) &= \mu_0 \int_V \mathbf{j}(t, \mathbf{x}') \delta(\mathbf{x} - \mathbf{x}') d^3x' \\ &\quad + \frac{\mu_0}{4\pi} \frac{\partial}{\partial t} \int_V \rho(t, \mathbf{x}') \nabla' \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right) d^3x' \\ &= \mu_0 \mathbf{j}(t, \mathbf{x}) + \mu_0 \frac{\partial}{\partial t} \epsilon_0 \mathbf{E}(t, \mathbf{x}) \end{aligned} \quad (1.22)$$

where, in the last step, we have assumed that a generalisation of equation (1.5) on page 4 to time-varying fields allows us to make the identification

$$\begin{aligned}
& \frac{1}{4\pi\epsilon_0} \frac{\partial}{\partial t} \int_V \rho(t, \mathbf{x}') \nabla' \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right) d^3x' \\
&= -\frac{1}{4\pi\epsilon_0} \frac{\partial}{\partial t} \int_V \rho(t, \mathbf{x}') \nabla \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right) d^3x' \\
&= \frac{\partial}{\partial t} \mathbf{E}(t, \mathbf{x})
\end{aligned} \tag{1.23}$$

Later, we will need to consider this formal result further. The result is Maxwell's source equation for the \mathbf{B} field

$$\nabla \times \mathbf{B}(t, \mathbf{x}) = \mu_0 \left(\mathbf{j}(t, \mathbf{x}) + \frac{\partial}{\partial t} \epsilon_0 \mathbf{E}(t, \mathbf{x}) \right) \tag{1.24}$$

where the last term $\partial \epsilon_0 \mathbf{E}(t, \mathbf{x}) / \partial t$ is the famous *displacement current*. This term was introduced, in a stroke of genius, by Maxwell in order to make the right hand side of this equation divergence free when $\mathbf{j}(t, \mathbf{x})$ is assumed to represent the density of the total electric current, which can be split up in “ordinary” conduction currents, polarisation currents and magnetisation currents. The displacement current is an extra term which behaves like a current density which flows in vacuum and, as we shall see later, its existence has very far-reaching physical consequences as it predicts the existence of electromagnetic radiation that can carry energy and momentum over very long distances, even in vacuum.

1.3.3 Electromotive force

If an electric field $\mathbf{E}(t, \mathbf{x})$, is applied to a conducting medium, a current density $\mathbf{j}(t, \mathbf{x})$ will be produced in this medium. There exist also hydrodynamical and chemical processes which can create currents. Under certain physical conditions, and for certain materials, one can sometimes assume a linear relationship between the current density \mathbf{j} and \mathbf{E} , called *Ohm's law*:

$$\mathbf{j}(t, \mathbf{x}) = \sigma \mathbf{E}(t, \mathbf{x}) \tag{1.25}$$

where σ is the *electric conductivity* (S/m). In the most general cases, for instance in an anisotropic conductor, σ is a tensor.

We can view Ohm's law, equation (1.25) above, as the first term in a Taylor expansion of the law $\mathbf{j}[\mathbf{E}(t, \mathbf{x})]$. This general law incorporates *non-linear effects* such as frequency mixing. Examples of media which are highly non-linear are semiconductors and plasma. We draw the attention to the fact that even in cases when the linear relation between \mathbf{E} and \mathbf{j} is a good approximation, we still have to use Ohm's law with care. The conductivity σ is, in general, time-dependent

(*temporal dispersive media*) but then it is often the case that equation (1.25) on the facing page is valid for each individual Fourier component of the field. We shall not, however, dwell upon such complicated cases here.

If the current is caused by an applied electric field $\mathbf{E}(t, \mathbf{x})$, this electric field will exert work on the charges in the medium and, unless the medium is superconducting, there will be some energy loss. The rate at which this energy is expended is $\mathbf{j} \cdot \mathbf{E}$ per unit volume. If \mathbf{E} is irrotational (conservative), \mathbf{j} will decay away with time. Stationary currents therefore require that an electric field which corresponds to an *electromotive force (EMF)* is present. In the presence of such a field \mathbf{E}^{EMF} , Ohm's law, equation (1.25) on the preceding page, takes the form

$$\mathbf{j} = \sigma(\mathbf{E}^{\text{stat}} + \mathbf{E}^{\text{EMF}}) \quad (1.26)$$

The electromotive force is defined as

$$\mathcal{E} = \oint_C (\mathbf{E}^{\text{stat}} + \mathbf{E}^{\text{EMF}}) \cdot d\mathbf{l} \quad (1.27)$$

where $d\mathbf{l}$ is a tangential line element of the closed loop C .

1.3.4 Faraday's law of induction

In subsection 1.1.2 we derived the differential equations for the electrostatic field. In particular, we derived equation (1.6) on page 4 which states that $\nabla \times \mathbf{E}^{\text{stat}}(\mathbf{x}) = \mathbf{0}$ and thus that \mathbf{E}^{stat} is a *conservative field* (it can be expressed as a gradient of a scalar field). This implies that the closed line integral of \mathbf{E}^{stat} in equation (1.27) above vanishes and that this equation becomes

$$\mathcal{E} = \oint_C \mathbf{E}^{\text{EMF}} \cdot d\mathbf{l} \quad (1.28)$$

It has been established experimentally that a nonconservative EMF field is produced in a closed circuit C if the magnetic flux through this circuit varies with time. This is formulated in *Faraday's law* which, in Maxwell's generalised form, reads

$$\begin{aligned} \mathcal{E}(t, \mathbf{x}) &= \oint_C \mathbf{E}(t, \mathbf{x}) \cdot d\mathbf{l} \\ &= -\frac{d}{dt} \Phi_m(t, \mathbf{x}) \\ &= -\frac{d}{dt} \int_S \mathbf{B}(t, \mathbf{x}) \cdot d\mathbf{S} \\ &= -\int_S d\mathbf{S} \cdot \frac{\partial}{\partial t} \mathbf{B}(t, \mathbf{x}) \end{aligned} \quad (1.29)$$

where Φ_m is the *magnetic flux* and S is the surface encircled by C which can be interpreted as a generic stationary "loop" and not necessarily as a conducting cir-

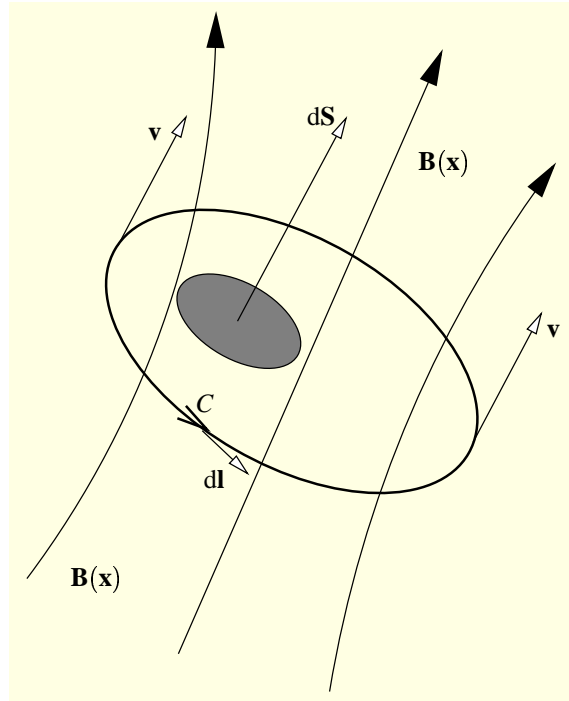


Figure 1.3. A loop C which moves with velocity \mathbf{v} in a spatially varying magnetic field $\mathbf{B}(\mathbf{x})$ will sense a varying magnetic flux during the motion.

cuit. Application of Stokes' theorem on this integral equation, transforms it into the differential equation

$$\nabla \times \mathbf{E}(t, \mathbf{x}) = -\frac{\partial}{\partial t} \mathbf{B}(t, \mathbf{x}) \quad (1.30)$$

which is valid for arbitrary variations in the fields and constitutes the Maxwell equation which explicitly connects electricity with magnetism.

Any change of the magnetic flux Φ_m will induce an EMF. Let us therefore consider the case, illustrated if figure 1.3, that the "loop" is moved in such a way that it links a magnetic field which varies during the movement. The *convective derivative* is evaluated according to the well-known operator formula

$$\frac{d}{dt} = \frac{\partial}{\partial t} + \mathbf{v} \cdot \nabla \quad (1.31)$$

which follows immediately from the rules of differentiation of an arbitrary differentiable function $f(t, \mathbf{x}(t))$. Applying this rule to Faraday's law, equation (1.29) on

page 11, we obtain

$$\begin{aligned}\mathcal{E}(t, \mathbf{x}) &= -\frac{d}{dt} \int_S \mathbf{B} \cdot d\mathbf{S} \\ &= -\int_S d\mathbf{S} \cdot \frac{\partial}{\partial t} \mathbf{B} - \int_S (\mathbf{v} \cdot \nabla) \mathbf{B} \cdot d\mathbf{S}\end{aligned}\quad (1.32)$$

In spatial differentiation \mathbf{v} is to be considered as constant, and equation (1.15) on page 7 holds also for time-varying fields:

$$\nabla \cdot \mathbf{B}(t, \mathbf{x}) = 0 \quad (1.33)$$

(it is one of Maxwell's equations) so that

$$\nabla \times (\mathbf{B} \times \mathbf{v}) = (\mathbf{v} \cdot \nabla) \mathbf{B} \quad (1.34)$$

allowing us to rewrite equation (1.32) above in the following way:

$$\begin{aligned}\mathcal{E}(t, \mathbf{x}) &= \oint_C \mathbf{E}^{\text{EMF}} \cdot d\mathbf{l} \\ &= -\frac{d}{dt} \int_S \mathbf{B} \cdot d\mathbf{S} \\ &= -\int_S \frac{\partial}{\partial t} \mathbf{B} \cdot d\mathbf{S} - \int_S \nabla \times (\mathbf{B} \times \mathbf{v}) \cdot d\mathbf{S}\end{aligned}\quad (1.35)$$

With Stokes' theorem applied to the last integral, we finally get

$$\begin{aligned}\mathcal{E}(t, \mathbf{x}) &= \oint_C \mathbf{E}^{\text{EMF}} \cdot d\mathbf{l} \\ &= -\int_S \frac{\partial}{\partial t} \mathbf{B} \cdot d\mathbf{S} - \oint_C (\mathbf{B} \times \mathbf{v}) \cdot d\mathbf{l}\end{aligned}\quad (1.36)$$

or, rearranging the terms,

$$\int_C (\mathbf{E}^{\text{EMF}} - \mathbf{v} \times \mathbf{B}) \cdot d\mathbf{l} = -\int_S \frac{\partial}{\partial t} \mathbf{B} \cdot d\mathbf{S} \quad (1.37)$$

where \mathbf{E}^{EMF} is the field which is induced in the "loop," *i.e.*, in the *moving* system.

The use of Stokes' theorem “backwards” on equation (1.37) yields

$$\nabla \times (\mathbf{E}^{\text{EMF}} - \mathbf{v} \times \mathbf{B}) = -\frac{\partial}{\partial t} \mathbf{B} \quad (1.38)$$

In the *fixed* system, an observer measures the electric field

$$\mathbf{E} = \mathbf{E}^{\text{EMF}} - \mathbf{v} \times \mathbf{B} \quad (1.39)$$

Hence, a moving observer measures the following *Lorentz force* on a charge q

$$q\mathbf{E}^{\text{EMF}} = q\mathbf{E} + q(\mathbf{v} \times \mathbf{B}) \quad (1.40)$$

corresponding to an “effective” electric field in the “loop” (moving observer)

$$\mathbf{E}^{\text{EMF}} = \mathbf{E} + (\mathbf{v} \times \mathbf{B}) \quad (1.41)$$

Hence, we can conclude that for a *stationary* observer, the Maxwell equation

$$\nabla \times \mathbf{E} = -\frac{\partial}{\partial t} \mathbf{B} \quad (1.42)$$

is indeed valid even if the “loop” is moving.

1.3.5 Maxwell's microscopic equations

We are now able to collect the results from the above considerations and formulate the equations of classical electrodynamics valid for arbitrary variations in time and space of the coupled electric and magnetic fields $\mathbf{E}(t, \mathbf{x})$ and $\mathbf{B}(t, \mathbf{x})$. The equations are

$$\nabla \cdot \mathbf{E} = \frac{\rho(t, \mathbf{x})}{\epsilon_0} \quad (1.43a)$$

$$\nabla \times \mathbf{E} + \frac{\partial \mathbf{B}}{\partial t} = \mathbf{0} \quad (1.43b)$$

$$\nabla \cdot \mathbf{B} = 0 \quad (1.43c)$$

$$\nabla \times \mathbf{B} - \epsilon_0 \mu_0 \frac{\partial \mathbf{E}}{\partial t} = \mu_0 \mathbf{j}(t, \mathbf{x}) \quad (1.43d)$$

In these equations $\rho(t, \mathbf{x})$ represents the total, possibly both time and space dependent, electric charge, *i.e.*, free as well as induced (polarisation) charges, and $\mathbf{j}(t, \mathbf{x})$ represents the total, possibly both time and space dependent, electric current, *i.e.*, conduction currents (motion of free charges) as well as all atomistic (polarisation,

magnetisation) currents. As they stand, the equations therefore incorporate the classical interaction between all electric charges and currents in the system and are called *Maxwell's microscopic equations*. Another name often used for them is the *Maxwell-Lorentz equations*. Together with the appropriate *constitutive relations*, which relate ρ and \mathbf{j} to the fields, and the initial and boundary conditions pertinent to the physical situation at hand, they form a system of well-posed partial differential equations which completely determine \mathbf{E} and \mathbf{B} .

1.3.6 Maxwell's macroscopic equations

The microscopic field equations (1.43) provide a correct classical picture for arbitrary field and source distributions, including both microscopic and macroscopic scales. However, for macroscopic substances it is sometimes convenient to introduce new derived fields which represent the electric and magnetic fields in which, in an average sense, the material properties of the substances are already included. These fields are the *electric displacement* \mathbf{D} and the *magnetising field* \mathbf{H} . In the most general case, these derived fields are complicated nonlocal, nonlinear functionals of the primary fields \mathbf{E} and \mathbf{B} :

$$\mathbf{D} = \mathbf{D}[t, \mathbf{x}; \mathbf{E}, \mathbf{B}] \quad (1.44a)$$

$$\mathbf{H} = \mathbf{H}[t, \mathbf{x}; \mathbf{E}, \mathbf{B}] \quad (1.44b)$$

Under certain conditions, for instance for very low field strengths, we may assume that the response of a substance is linear so that

$$\mathbf{D} = \varepsilon(t, \mathbf{x})\mathbf{E} \quad (1.45)$$

$$\mathbf{H} = \mu^{-1}(t, \mathbf{x})\mathbf{B} \quad (1.46)$$

i.e., that the derived fields are linearly proportional to the primary fields and that the electric displacement (magnetising field) is only dependent on the electric (magnetic) field.

The field equations expressed in terms of the derived field quantities \mathbf{D} and \mathbf{H} are

$$\nabla \cdot \mathbf{D} = \rho(t, \mathbf{x}) \quad (1.47a)$$

$$\nabla \times \mathbf{E} + \frac{\partial \mathbf{B}}{\partial t} = \mathbf{0} \quad (1.47b)$$

$$\nabla \cdot \mathbf{B} = 0 \quad (1.47c)$$

$$\nabla \times \mathbf{H} - \frac{\partial \mathbf{D}}{\partial t} = \mathbf{j}(t, \mathbf{x}) \quad (1.47d)$$

and are called *Maxwell's macroscopic equations*.

1.4 Electromagnetic Duality

If we look more closely at the microscopic Maxwell equations (1.48), we see that they exhibit a certain, albeit not a complete, symmetry. Let us for explicitness denote the electric charge density $\rho = \rho(t, \mathbf{x})$ by ρ_e and the electric current density $\mathbf{j} = \mathbf{j}(t, \mathbf{x})$ by \mathbf{j}_e . We further make the *ad hoc* assumption that there exist *magnetic monopoles* represented by a *magnetic charge density*, denoted $\rho_m = \rho_m(t, \mathbf{x})$, and a *magnetic current density*, denoted $\mathbf{j}_m = \mathbf{j}_m(t, \mathbf{x})$. With these new quantities included in the theory, the Maxwell equations can be written

$$\nabla \cdot \mathbf{E} = \frac{\rho_e}{\epsilon_0} \quad (1.48a)$$

$$\nabla \times \mathbf{E} + \frac{\partial \mathbf{B}}{\partial t} = -\mu_0 \mathbf{j}_m \quad (1.48b)$$

$$\nabla \cdot \mathbf{B} = \mu_0 \rho_m \quad (1.48c)$$

$$\nabla \times \mathbf{B} - \frac{1}{c^2} \frac{\partial \mathbf{E}}{\partial t} = \mu_0 \mathbf{j}_e \quad (1.48d)$$

We shall call these equations the *Dirac-Maxwell equations* or the *electromagnetodynamic equations*

Taking the divergence of (1.48b), we find that

$$\nabla \cdot (\nabla \times \mathbf{E}) = -\frac{\partial}{\partial t} (\nabla \cdot \mathbf{B}) - \mu_0 \nabla \cdot \mathbf{j}_m \equiv 0 \quad (1.49)$$

where we used the fact that, according to formula (M.79) on page 189, the divergence of a curl always vanishes. Using (1.48c) to rewrite this relation, we obtain the *equation of continuity for magnetic monopoles*

$$\frac{\partial \rho_m}{\partial t} + \nabla \cdot \mathbf{j}_m = 0 \quad (1.50)$$

which has the same form as that for the electric monopoles (electric charges) and currents, equation (1.21) on page 9.

We notice that the new equations (1.48) above exhibit the following symmetry (recall that $\epsilon_0 \mu_0 = 1/c^2$):

$$\mathbf{E} \rightarrow c\mathbf{B} \quad (1.51a)$$

$$\mathbf{B} \rightarrow -\frac{1}{c}\mathbf{E} \quad (1.51b)$$

$$\rho_e \rightarrow \frac{1}{c}\rho_m \quad (1.51c)$$

$$\rho_m \rightarrow -c\rho_e \quad (1.51d)$$

$$\mathbf{j}_e \rightarrow \frac{1}{c}\mathbf{j}_m \quad (1.51e)$$

$$\mathbf{j}_m \rightarrow -c\mathbf{j}_e \quad (1.51f)$$

which is a particular case ($\theta = \pi/2$) of the general *duality transformation* (depicted by the *Hodge star operator*)

$$*\mathbf{E} = \mathbf{E} \cos \theta + c\mathbf{B} \sin \theta \quad (1.52a)$$

$$*\mathbf{B} = -\frac{1}{c}\mathbf{E} \sin \theta + \mathbf{B} \cos \theta \quad (1.52b)$$

$$*\rho_e = \rho_e \cos \theta + \frac{1}{c}\rho_m \sin \theta \quad (1.52c)$$

$$*\rho_m = -c\rho_e \sin \theta + \rho_m \cos \theta \quad (1.52d)$$

$$*\mathbf{j}_e = \mathbf{j}_e \cos \theta + \frac{1}{c}\mathbf{j}_m \sin \theta \quad (1.52e)$$

$$*\mathbf{j}_m = -c\mathbf{j}_e \sin \theta + \mathbf{j}_m \cos \theta \quad (1.52f)$$

which leaves the Dirac-Maxwell equations, and hence the physics they describe (often referred to as *electromagnetodynamics*), invariant. Since \mathbf{E} and \mathbf{j}_e are (true or polar) vectors, \mathbf{B} a pseudovector (axial vector), ρ_e a (true) scalar, then ρ_m and θ , which behaves as a *mixing angle* in a two-dimensional “charge space,” must be pseudoscalars and \mathbf{j}_m a pseudovector.

▷ DUALITY OF THE ELECTROMAGNETODYNAMIC EQUATIONS

EXAMPLE 1.1

Show that the symmetric, electromagnetodynamic form of Maxwell’s equations (the Dirac-Maxwell equations), equations (1.48) on the facing page are invariant under the duality transformation (1.52).

Explicit application of the transformation yields

$$\begin{aligned}
\nabla \cdot {}^* \mathbf{E} &= \nabla \cdot (\mathbf{E} \cos \theta + c \mathbf{B} \sin \theta) = \frac{\rho_e}{\epsilon_0} \cos \theta + c \mu_0 \rho_m \sin \theta \\
&= \frac{1}{\epsilon_0} \left(\rho_e \cos \theta + \frac{1}{c} \rho_m \sin \theta \right) = \frac{{}^* \rho_e}{\epsilon_0} \quad (1.53) \\
\nabla \times {}^* \mathbf{E} + \frac{\partial {}^* \mathbf{B}}{\partial t} &= \nabla \times (\mathbf{E} \cos \theta + c \mathbf{B} \sin \theta) + \frac{\partial}{\partial t} \left(-\frac{1}{c} \mathbf{E} \sin \theta + \mathbf{B} \cos \theta \right) \\
&= -\frac{\partial \mathbf{B}}{\partial t} \cos \theta - \mu_0 \mathbf{j}_m \cos \theta + \frac{1}{c} \frac{\partial \mathbf{E}}{\partial t} \sin \theta + c \mu_0 \mathbf{j}_e \sin \theta \\
&\quad - \frac{1}{c} \frac{\partial \mathbf{E}}{\partial t} \sin \theta + \frac{\partial \mathbf{B}}{\partial t} \cos \theta = -\mu_0 \mathbf{j}_m \cos \theta + c \mu_0 \mathbf{j}_e \sin \theta \\
&= -\mu_0 (-c \mathbf{j}_e \sin \theta + \mathbf{j}_m \cos \theta) = -\mu_0 {}^* \mathbf{j}_m \quad (1.54)
\end{aligned}$$

and analogously for the other two Dirac-Maxwell equations. QED ■

—END OF EXAMPLE 1.1◀

The invariance of the Dirac-Maxwell equations under the similarity transformation means that the amount of magnetic monopole density ρ_m is irrelevant for the physics as long as the ratio $\rho_m/\rho_e = \tan \theta$ is kept constant. So whether we assume that the particles are only electrically charged or have also a magnetic charge with a given, fixed ratio between the two types of charges is a matter of convention, as long as we assume that this fraction is *the same for all particles*. By varying the mixing angle θ we can change the fraction of magnetic monopoles at will without changing the laws of electrodynamics. For $\theta = 0$ we recover the usual Maxwell electrodynamics as we know it.

EXAMPLE 1.2

▷ MAXWELL FROM DIRAC-MAXWELL EQUATIONS FOR A FIXED MIXING ANGLE

Show that for a fixed mixing angle θ such that

$$\rho_m = c \rho_e \tan \theta \quad (1.55a)$$

$$\mathbf{j}_m = c \mathbf{j}_e \tan \theta \quad (1.55b)$$

the Dirac-Maxwell equations reduce to the usual Maxwell equations.

Explicit application of the fixed mixing angle conditions on the duality transformation (1.52) on the previous page yields

$$\begin{aligned}
{}^* \rho_e &= \rho_e \cos \theta + \frac{1}{c} \rho_m \sin \theta = \rho_e \cos \theta + \frac{1}{c} c \rho_e \tan \theta \sin \theta \\
&= \frac{1}{\cos \theta} (\rho_e \cos^2 \theta + \rho_e \sin^2 \theta) = \frac{1}{\cos \theta} \rho_e \quad (1.56a)
\end{aligned}$$

$${}^* \rho_m = -c \rho_e \sin \theta + c \rho_e \tan \theta \cos \theta = -c \rho_e \sin \theta + c \rho_e \sin \theta = 0 \quad (1.56b)$$

$${}^* \mathbf{j}_e = \mathbf{j}_e \cos \theta + \mathbf{j}_e \tan \theta \sin \theta = \frac{1}{\cos \theta} (\mathbf{j}_e \cos^2 \theta + \mathbf{j}_e \sin^2 \theta) = \frac{1}{\cos \theta} \mathbf{j}_e \quad (1.56c)$$

$${}^* \mathbf{j}_m = -c \mathbf{j}_e \sin \theta + c \mathbf{j}_e \tan \theta \cos \theta = -c \mathbf{j}_e \sin \theta + c \mathbf{j}_e \sin \theta = 0 \quad (1.56d)$$

Hence, a fixed mixing angle, or, equivalently, a fixed ratio between the electric and magnetic charges/currents, “hides” the magnetic monopole influence (ρ_m and \mathbf{j}_m) on the dynamic equations. Furthermore, we notice that

$$\begin{aligned}
 \nabla \cdot \mathbf{E} &= \nabla \cdot \mathbf{E} \cos \theta + c \nabla \cdot \mathbf{B} \sin \theta = \nabla \cdot \mathbf{E} \cos \theta + c \mu_0 \rho_m \sin \theta \\
 &= \nabla \cdot \mathbf{E} \cos \theta + c^2 \mu_0 \rho_e \tan \theta \sin \theta = \nabla \cdot \mathbf{E} \cos \theta + \frac{\rho_e}{\epsilon_0} \tan \theta \sin \theta \\
 &= \nabla \cdot \mathbf{E} \cos \theta + \frac{\rho_e \sin^2 \theta}{\epsilon_0 \cos \theta} = \frac{1}{\cos \theta} \frac{\rho_e}{\epsilon_0}
 \end{aligned} \tag{1.57}$$

or

$$\nabla \cdot \mathbf{E} = \frac{1}{\cos^2 \theta} \frac{\rho_e}{\epsilon_0} (1 - \sin^2 \theta) = \frac{\rho_e}{\epsilon_0} \tag{1.58}$$

and so on for the other equations.

QED ■

—END OF EXAMPLE 1.2◀

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Electromagnetic Waves

Maxwell's microscopic equations (1.43) on page 14, which are usually written in the following form

$$\nabla \cdot \mathbf{E} = \frac{\rho(t, \mathbf{x})}{\epsilon_0} \quad (2.1a)$$

$$\nabla \times \mathbf{E} = -\frac{\partial}{\partial t} \mathbf{B} \quad (2.1b)$$

$$\nabla \cdot \mathbf{B} = 0 \quad (2.1c)$$

$$\nabla \times \mathbf{B} = \mu_0 \mathbf{j}(t, \mathbf{x}) + \epsilon_0 \mu_0 \frac{\partial}{\partial t} \mathbf{E} \quad (2.1d)$$

can be viewed as an axiomatic basis for classical electrodynamics. In particular, these equations are well suited for calculating the electric and magnetic fields \mathbf{E} and \mathbf{B} from given, prescribed charge distributions $\rho(t, \mathbf{x})$ and current distributions $\mathbf{j}(t, \mathbf{x})$ of arbitrary time- and space-dependent form.

However, as is well known from the theory of differential equations, these four first order, coupled partial differential vector equations can be rewritten as two uncoupled, second order partial equations, one for \mathbf{E} and one for \mathbf{B} . We shall derive the second order equation for \mathbf{E} , which, as we shall see is a *homogeneous wave equation*, and then discuss the implications of this equation. We shall also show how the \mathbf{B} field can be easily calculated from the solution of the \mathbf{E} equation.

2.1 The wave equation

Let us consider a volume with no net charge, $\rho = 0$, and no electromotive force $\mathbf{E}^{\text{EMF}} = \mathbf{0}$. Taking the curl of (2.1b) and using (2.1d), we obtain

$$\begin{aligned}\nabla \times (\nabla \times \mathbf{E}) &= -\frac{\partial}{\partial t}(\nabla \times \mathbf{B}) \\ &= -\mu_0 \frac{\partial}{\partial t} \left(\mathbf{j} + \varepsilon_0 \frac{\partial \mathbf{E}}{\partial t} \right)\end{aligned}\quad (2.2)$$

According to the operator triple product “bac-cab” rule equation (F.69) on page 168

$$\nabla \times (\nabla \times \mathbf{E}) = \nabla(\nabla \cdot \mathbf{E}) - \nabla^2 \mathbf{E} \quad (2.3)$$

Furthermore, since $\rho = 0$, equation (2.1a) on the previous page yields

$$\nabla \cdot \mathbf{E} = 0 \quad (2.4)$$

and since $\mathbf{E}^{\text{EMF}} = \mathbf{0}$, Ohm’s law, equation (1.26) on page 11, yields

$$\mathbf{j} = \sigma \mathbf{E} \quad (2.5)$$

we find that Equation (2.2) can be rewritten

$$\nabla^2 \mathbf{E} - \mu_0 \frac{\partial}{\partial t} \left(\sigma \mathbf{E} + \varepsilon_0 \frac{\partial \mathbf{E}}{\partial t} \right) = \mathbf{0} \quad (2.6)$$

or, also using equation (1.9) on page 5,

$$\nabla^2 \mathbf{E} - \mu_0 \sigma \frac{\partial \mathbf{E}}{\partial t} - \frac{1}{c^2} \frac{\partial^2 \mathbf{E}}{\partial t^2} = \mathbf{0} \quad (2.7)$$

which is the *homogeneous wave equation* for \mathbf{E} .

We look for a solution in the form of a *time-harmonic wave*, and make therefore the following *Fourier component Ansatz*

$$\mathbf{E} = \mathbf{E}_0(\mathbf{x})e^{-i\omega t} \quad (2.8)$$

Inserting this into equation (2.7), we obtain

$$\begin{aligned}\nabla^2 \mathbf{E} - \mu_0 \sigma \frac{\partial}{\partial t} \mathbf{E}_0(\mathbf{x})e^{-i\omega t} - \frac{1}{c^2} \frac{\partial^2}{\partial t^2} \mathbf{E}_0(\mathbf{x})e^{-i\omega t} \\ = \nabla^2 \mathbf{E} - \mu_0 \sigma (-i\omega) \mathbf{E}_0(\mathbf{x})e^{-i\omega t} - \frac{1}{c^2} (-i\omega)^2 \mathbf{E}_0(\mathbf{x})e^{-i\omega t} \\ = \nabla^2 \mathbf{E} - \mu_0 \sigma (-i\omega) \mathbf{E} - \frac{1}{c^2} (-i\omega)^2 \mathbf{E} = \mathbf{0}\end{aligned}\quad (2.9)$$

which we can rewrite as

$$\nabla^2 \mathbf{E} + \frac{\omega^2}{c^2} \left(1 + i \frac{\sigma}{\varepsilon_0 \omega} \right) \mathbf{E} = \nabla^2 \mathbf{E} + \frac{\omega^2}{c^2} \left(1 + \frac{i}{\tau \omega} \right) \mathbf{E} = \mathbf{0} \quad (2.10)$$

The quantity $\tau = \varepsilon_0/\sigma$ is called the *relaxation time* of the medium in question. In

the limit of long τ , Equation (2.10) tends to

$$\nabla^2 \mathbf{E} + \frac{\omega^2}{c^2} \mathbf{E} = \mathbf{0} \quad (2.11)$$

which is a *time-independent wave equation* for \mathbf{E} , representing weakly damped propagating waves. In the short τ limit we have instead

$$\nabla^2 \mathbf{E} + i\omega\mu_0\sigma\mathbf{E} = \mathbf{0} \quad (2.12)$$

which is a *time-independent diffusion equation* for \mathbf{E} .

For most metals $\tau \sim 10^{-14}$ s, which means that the diffusion picture is good for all frequencies lower than optical frequencies. Hence, in metallic conductors, the propagation term $\partial^2 \mathbf{E}/c^2 \partial t^2$ is negligible even for VHF, UHF, and SHF signals. Alternatively, we may say that the displacement current $\epsilon_0 \partial \mathbf{E}/\partial t$ is negligible relative to the conduction current $\mathbf{j} = \sigma \mathbf{E}$.

If we introduce the *vacuum wave number*

$$k = \frac{\omega}{c} \quad (2.13)$$

we can write, using the fact that $c = 1/\sqrt{\epsilon_0\mu_0}$ according to equation (1.9) on page 5,

$$\frac{1}{\tau\omega} = \frac{\sigma}{\epsilon_0\omega} = \frac{\sigma}{\epsilon_0} \frac{1}{ck} = \frac{\sigma}{k} \sqrt{\frac{\mu_0}{\epsilon_0}} = \frac{\sigma}{k} R_0 \quad (2.14)$$

where in the last step we introduced the *characteristic impedance* for vacuum

$$R_0 = \sqrt{\frac{\mu_0}{\epsilon_0}} \approx 376.7 \Omega \quad (2.15)$$

2.2 Plane waves

Consider now the case where all fields depend only on the distance ζ to a given plane with unit normal $\hat{\mathbf{n}}$. Then the *del* operator becomes

$$\nabla = \hat{\mathbf{n}} \frac{\partial}{\partial \zeta} \quad (2.16)$$

and Maxwell's equations attain the form

$$\hat{\mathbf{n}} \cdot \frac{\partial \mathbf{E}}{\partial \zeta} = 0 \quad (2.17a)$$

$$\hat{\mathbf{n}} \times \frac{\partial \mathbf{E}}{\partial \zeta} = -\frac{\partial \mathbf{B}}{\partial t} \quad (2.17b)$$

$$\hat{\mathbf{n}} \cdot \frac{\partial \mathbf{B}}{\partial \zeta} = 0 \quad (2.17c)$$

$$\hat{\mathbf{n}} \times \frac{\partial \mathbf{B}}{\partial \zeta} = \mu_0 \mathbf{j}(t, \mathbf{x}) + \varepsilon_0 \mu_0 \frac{\partial \mathbf{E}}{\partial t} = \mu_0 \sigma \mathbf{E} + \varepsilon_0 \mu_0 \frac{\partial \mathbf{E}}{\partial t} \quad (2.17d)$$

Scalar multiplying (2.17d) by $\hat{\mathbf{n}}$, we find that

$$0 = \hat{\mathbf{n}} \cdot \left(\hat{\mathbf{n}} \times \frac{\partial \mathbf{B}}{\partial \zeta} \right) = \hat{\mathbf{n}} \cdot \left(\mu_0 \sigma + \varepsilon_0 \mu_0 \frac{\partial}{\partial t} \right) \mathbf{E} \quad (2.18)$$

which simplifies to the first-order ordinary differential equation for the normal component E_n of the electric field

$$\frac{dE_n}{dt} + \frac{\sigma}{\varepsilon_0} E_n = 0 \quad (2.19)$$

with the solution

$$E_n = E_{n_0} e^{-\sigma t / \varepsilon_0} = E_{n_0} e^{-t / \tau} \quad (2.20)$$

This, together with (2.17a), shows that the *longitudinal component* of \mathbf{E} , *i.e.*, the component which is perpendicular to the plane surface is independent of ζ and has a time dependence which exhibits an exponential decay, with a decrement given by the relaxation time τ in the medium.

Scalar multiplying (2.17b) by $\hat{\mathbf{n}}$, we similarly find that

$$0 = \hat{\mathbf{n}} \cdot \left(\hat{\mathbf{n}} \times \frac{\partial \mathbf{E}}{\partial \zeta} \right) = -\hat{\mathbf{n}} \cdot \frac{\partial \mathbf{B}}{\partial t} \quad (2.21)$$

or

$$\hat{\mathbf{n}} \cdot \frac{\partial \mathbf{B}}{\partial t} = 0 \quad (2.22)$$

From this, and (2.17c), we conclude that the only longitudinal component of \mathbf{B} must be constant in both time and space. In other words, the only non-static solution must consist of *transverse components*.

2.2.1 Telegrapher's equation

In analogy with equation (2.7) on page 24, we can easily derive the equation

$$\frac{\partial^2 \mathbf{E}}{\partial \zeta^2} - \mu_0 \sigma \frac{\partial \mathbf{E}}{\partial t} - \frac{1}{c^2} \frac{\partial^2 \mathbf{E}}{\partial t^2} = \mathbf{0} \quad (2.23)$$

This equation, which describes the propagation of plane waves in a conducting medium, is called the *telegrapher's equation*. If the medium is an insulator so that $\sigma = 0$, then the equation takes the form of the *one-dimensional wave equation*

$$\frac{\partial^2 \mathbf{E}}{\partial \zeta^2} - \frac{1}{c^2} \frac{\partial^2 \mathbf{E}}{\partial t^2} = \mathbf{0} \quad (2.24)$$

As is well known, each component of this equation has a solution which can be written

$$E_i = f(\zeta - ct) + g(\zeta + ct), \quad i = 1, 2, 3 \quad (2.25)$$

where f and g are arbitrary (non-pathological) functions of their respective arguments. This general solution represents perturbations which propagate along ζ , where the f perturbation propagates in the positive ζ direction and the g perturbation propagates in the negative ζ direction.

If we assume that \mathbf{E} is time-harmonic, *i.e.*, can be represented by a Fourier component proportional to $\exp -i\omega t$, the solution of Equation (2.24) becomes

$$\mathbf{E} = \mathbf{E}_0 e^{-i(\omega t \pm k \zeta)} \quad (2.26)$$

By introducing the *wave vector*

$$\mathbf{k} = k \hat{\mathbf{n}} = \frac{\omega}{c} \hat{\mathbf{n}} = \frac{\omega}{c} \hat{\mathbf{k}} \quad (2.27)$$

this solution can be written as

$$\mathbf{E} = \mathbf{E}_0 e^{i(\mathbf{k} \cdot \mathbf{x} - \omega t)} \quad (2.28)$$

Let us consider the minus sign in the exponent in equation (2.26) above. This corresponds to a wave which propagates in the direction of increasing ζ . Inserting

this solution into equation (2.17b) on page 26, gives

$$\hat{\mathbf{n}} \times \frac{\partial \mathbf{E}}{\partial \zeta} = i\omega \mathbf{B} = ik\hat{\mathbf{n}} \times \mathbf{E} \quad (2.29)$$

or, solving for \mathbf{B} ,

$$\mathbf{B} = \frac{k}{\omega} \hat{\mathbf{n}} \times \mathbf{E} = \frac{1}{\omega} \mathbf{k} \times \mathbf{E} = \frac{1}{c} \hat{\mathbf{k}} \times \mathbf{E} = \sqrt{\epsilon_0 \mu_0} \hat{\mathbf{n}} \times \mathbf{E} \quad (2.30)$$

Hence, to each transverse component of \mathbf{E} , there exists an associated magnetic field given by equation (2.30) above. If \mathbf{E} and/or \mathbf{B} has a direction in space which is constant in time, we have a *plane polarised wave* (or *linearly polarised wave*).

2.2.2 Waves in conductive media

Assuming that our medium has a finite conductivity σ , and making the time-harmonic wave Ansatz in equation (2.23) on the preceding page, we find that the *time-independent telegrapher's equation* can be written

$$\frac{\partial^2 \mathbf{E}}{\partial \zeta^2} + \epsilon_0 \mu_0 \omega^2 \mathbf{E} + i\mu_0 \sigma \omega \mathbf{E} = \frac{\partial^2 \mathbf{E}}{\partial \zeta^2} + K^2 \mathbf{E} = \mathbf{0} \quad (2.31)$$

where

$$K^2 = \epsilon_0 \mu_0 \omega^2 \left(1 + i \frac{\sigma}{\epsilon_0 \omega} \right) = \frac{\omega^2}{c^2} \left(1 + i \frac{\sigma}{\epsilon_0 \omega} \right) = k^2 \left(1 + i \frac{\sigma}{\epsilon_0 \omega} \right) \quad (2.32)$$

where, in the last step, equation (2.13) on page 25 was used to introduce the wave number k . Taking the square root of this expression, we obtain

$$K = k \sqrt{1 + i \frac{\sigma}{\epsilon_0 \omega}} = \alpha + i\beta \quad (2.33)$$

Squaring, we find that

$$k^2 \left(1 + i \frac{\sigma}{\epsilon_0 \omega} \right) = (\alpha^2 - \beta^2) + 2i\alpha\beta \quad (2.34)$$

or

$$\beta^2 = \alpha^2 - k^2 \quad (2.35)$$

$$\alpha\beta = \frac{k^2 \sigma}{2\epsilon_0 \omega} \quad (2.36)$$

Squaring the latter and combining with the former, we obtain the second order algebraic equation (in α^2)

$$\alpha^2(\alpha^2 - k^2) = \frac{k^4 \sigma^2}{4\varepsilon_0^2 \omega^2} \quad (2.37)$$

which can be easily solved and one finds that

$$\alpha = k \sqrt{\frac{\sqrt{1 + \left(\frac{\sigma}{\varepsilon_0 \omega}\right)^2} + 1}{2}} \quad (2.38a)$$

$$\beta = k \sqrt{\frac{\sqrt{1 + \left(\frac{\sigma}{\varepsilon_0 \omega}\right)^2} - 1}{2}} \quad (2.38b)$$

Hence, the solution of the time-independent telegrapher's equation, equation (2.31) on the facing page, can be written

$$\mathbf{E} = \mathbf{E}_0 e^{-\beta \zeta} e^{i(\alpha \zeta - \omega t)} \quad (2.39)$$

With the aid of equation (2.30) on the preceding page we can calculate the associated magnetic field, and find that it is given by

$$\mathbf{B} = \frac{1}{\omega} K \hat{\mathbf{k}} \times \mathbf{E} = \frac{1}{\omega} (\hat{\mathbf{k}} \times \mathbf{E})(\alpha + i\beta) = \frac{1}{\omega} (\hat{\mathbf{k}} \times \mathbf{E}) |A| e^{i\gamma} \quad (2.40)$$

where we have, in the last step, rewritten $\alpha + i\beta$ in the amplitude-phase form $|A| \exp\{i\gamma\}$. From the above, we immediately see that \mathbf{E} is damped and that \mathbf{E} and \mathbf{B} in the wave are out of phase.

In the case that $\varepsilon_0 \omega \ll \sigma$, we can approximate K as follows:

$$\begin{aligned} K &= k \left(1 + i \frac{\sigma}{\varepsilon_0 \omega}\right)^{\frac{1}{2}} = k \left[i \frac{\sigma}{\varepsilon_0 \omega} \left(1 - i \frac{\varepsilon_0 \omega}{\sigma}\right)\right]^{\frac{1}{2}} \approx k(1 + i) \sqrt{\frac{\sigma}{2\varepsilon_0 \omega}} \\ &= \sqrt{\varepsilon_0 \mu_0} \omega (1 + i) \sqrt{\frac{\sigma}{2\varepsilon_0 \omega}} = (1 + i) \sqrt{\frac{\mu_0 \sigma \omega}{2}} \end{aligned} \quad (2.41)$$

From this analysis we conclude that when the wave impinges perpendicularly upon the medium, the fields are given, *inside* this medium, by

$$\mathbf{E}' = \mathbf{E}_0 \exp\left\{-\sqrt{\frac{\mu_0 \sigma \omega}{2}} \zeta\right\} \exp\left\{i\left(\sqrt{\frac{\mu_0 \sigma \omega}{2}} \zeta - \omega t\right)\right\} \quad (2.42a)$$

$$\mathbf{B}' = (1 + i) \sqrt{\frac{\mu_0 \sigma}{2\omega}} (\hat{\mathbf{n}} \times \mathbf{E}') \quad (2.42b)$$

Hence, both fields fall off by a factor $1/e$ at a distance

$$\delta = \sqrt{\frac{2}{\mu_0 \sigma \omega}} \quad (2.43)$$

This distance δ is called the *skin depth*.

2.3 Observables and averages

In the above we have used *complex notation* quite extensively. This is for mathematical convenience only. For instance, in this notation differentiations are almost trivial to perform. However, every *physical measurable* quantity is always real valued. *I.e.*, “ $\mathbf{E}_{\text{physical}} = \text{Re} \{ \mathbf{E}_{\text{mathematical}} \}$.” It is particularly important to remember this when one works with products of physical quantities.

Generally speaking, we tend to measure temporal averages ($\langle \rangle$) of our physical observables. If we have two physical vectors \mathbf{F} and \mathbf{G} which both are time harmonic, *i.e.*, can be represented by Fourier components proportional to $\exp\{-i\omega t\}$, it is easy to show that the average of the product of the two physical quantities represented by \mathbf{F} and \mathbf{G} can be expressed as

$$\langle \text{Re} \{ \mathbf{F} \} \cdot \text{Re} \{ \mathbf{G} \} \rangle = \frac{1}{2} \text{Re} \{ \mathbf{F} \cdot \mathbf{G}^* \} = \frac{1}{2} \text{Re} \{ \mathbf{F}^* \cdot \mathbf{G} \} \quad (2.44)$$

where $*$ denotes complex conjugate.

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Electromagnetic Potentials

Instead of expressing the laws of electrodynamics in terms of electric and magnetic fields, it turns out that it is often more convenient to express the theory in terms of potentials. In this Chapter we will introduce and study the properties of such potentials.

3.1 The electrostatic scalar potential

As we saw in equation (1.6) on page 4, the electrostatic field $\mathbf{E}^{\text{stat}}(\mathbf{x})$ is irrotational. Hence, it may be expressed in terms of the gradient of a scalar field. If we denote this scalar field by $-\phi^{\text{stat}}(\mathbf{x})$, we get

$$\mathbf{E}^{\text{stat}}(\mathbf{x}) = -\nabla\phi^{\text{stat}}(\mathbf{x}) \quad (3.1)$$

Taking the divergence of this and using equation (1.7) on page 4, we obtain *Poisson's equation*

$$\nabla^2\phi^{\text{stat}}(\mathbf{x}) = -\nabla\cdot\mathbf{E}^{\text{stat}}(\mathbf{x}) = -\frac{\rho(\mathbf{x})}{\epsilon_0} \quad (3.2)$$

A comparison with the definition of \mathbf{E}^{stat} , namely equation (1.5) on page 4, after the ∇ has been moved out of the integral, shows that this equation has the solution

$$\phi^{\text{stat}}(\mathbf{x}) = \frac{1}{4\pi\epsilon_0} \int_V \frac{\rho(\mathbf{x}')}{|\mathbf{x} - \mathbf{x}'|} d^3x' + \alpha \quad (3.3)$$

where the integration is taken over all source points \mathbf{x}' at which the charge density $\rho(\mathbf{x}')$ is non-zero and α is an arbitrary quantity which has a vanishing gradient. An example of such a quantity is a scalar constant. The scalar function $\phi^{\text{stat}}(\mathbf{x})$ in equation (3.3) above is called the *electrostatic scalar potential*.

3.2 The magnetostatic vector potential

Consider the equations of magnetostatics (1.20) on page 8. From equation (M.79) on page 189 we know that any 3D vector \mathbf{a} has the property that $\nabla \cdot (\nabla \times \mathbf{a}) \equiv 0$ and in the derivation of equation (1.15) on page 7 in magnetostatics we found that $\nabla \cdot \mathbf{B}^{\text{stat}}(\mathbf{x}) = 0$. We therefore realise that we can always write

$$\mathbf{B}^{\text{stat}}(\mathbf{x}) = \nabla \times \mathbf{A}^{\text{stat}}(\mathbf{x}) \quad (3.4)$$

where $\mathbf{A}^{\text{stat}}(\mathbf{x})$ is called the *magnetostatic vector potential*.

We saw above that the electrostatic potential (as any scalar potential) is not unique: we may, without changing the physics, add to it a quantity whose spatial gradient vanishes. A similar arbitrariness is true also for the magnetostatic vector potential.

In the magnetostatic case, we may start from Biot-Savart's law as expressed by equation (1.13) on page 7 and "move the ∇ out of the integral:"

$$\begin{aligned} \mathbf{B}^{\text{stat}}(\mathbf{x}) &= \frac{\mu_0}{4\pi} \int_V \mathbf{j}(\mathbf{x}') \times \frac{\mathbf{x} - \mathbf{x}'}{|\mathbf{x} - \mathbf{x}'|^3} d^3x' \\ &= -\frac{\mu_0}{4\pi} \int_V \mathbf{j}(\mathbf{x}') \times \nabla \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right) d^3x' \\ &= \nabla \times \frac{\mu_0}{4\pi} \int_V \frac{\mathbf{j}(\mathbf{x}')}{|\mathbf{x} - \mathbf{x}'|} d^3x' \end{aligned} \quad (3.5)$$

An identification of terms allows us to define the static vector potential as

$$\mathbf{A}^{\text{stat}}(\mathbf{x}) = \frac{\mu_0}{4\pi} \int_V \frac{\mathbf{j}(\mathbf{x}')}{|\mathbf{x} - \mathbf{x}'|} d^3x' + \mathbf{a}(\mathbf{x}) \quad (3.6)$$

where $\mathbf{a}(\mathbf{x})$ is an arbitrary vector field whose curl vanishes. From equation (M.75) on page 188 we know that such a vector can always be written as the gradient of a scalar field.

3.3 The electromagnetic scalar and vector potentials

Let us now generalise the static analysis above to the electrodynamic case, *i.e.*, the case with temporal and spatial dependent sources $\rho(t, \mathbf{x})$ and $\mathbf{j}(t, \mathbf{x})$, and corresponding fields $\mathbf{E}(t, \mathbf{x})$ and $\mathbf{B}(t, \mathbf{x})$, as described by Maxwell's equations (1.43)

on page 14. In other words, let us study the *electromagnetic potentials* $\phi(t, \mathbf{x})$ and $\mathbf{A}(t, \mathbf{x})$.

From equation (1.43c) on page 14 we note that also in electrodynamics the homogeneous equation $\nabla \cdot \mathbf{B}(t, \mathbf{x}) = 0$ remains valid. Because of this divergence-free nature of the time- and space-dependent magnetic field, we can express it as the curl of an *electromagnetic vector potential*:

$$\mathbf{B}(t, \mathbf{x}) = \nabla \times \mathbf{A}(t, \mathbf{x}) \quad (3.7)$$

Inserting this expression into the other homogeneous Maxwell equation, equation (1.30) on page 12, we obtain

$$\nabla \times \mathbf{E}(t, \mathbf{x}) = -\frac{\partial}{\partial t} [\nabla \times \mathbf{A}(t, \mathbf{x})] = -\nabla \times \frac{\partial}{\partial t} \mathbf{A}(t, \mathbf{x}) \quad (3.8)$$

or, rearranging the terms,

$$\nabla \times \left(\mathbf{E}(t, \mathbf{x}) + \frac{\partial}{\partial t} \mathbf{A}(t, \mathbf{x}) \right) = \mathbf{0} \quad (3.9)$$

As before we utilise the vanishing curl of a vector expression to write this vector expression as the gradient of a scalar function. If, in analogy with the electrostatic case, we introduce the *electromagnetic scalar potential* function $-\phi(t, \mathbf{x})$, equation (3.9) becomes equivalent to

$$\mathbf{E}(t, \mathbf{x}) + \frac{\partial}{\partial t} \mathbf{A}(t, \mathbf{x}) = -\nabla \phi(t, \mathbf{x}) \quad (3.10)$$

This means that in electrodynamics, $\mathbf{E}(t, \mathbf{x})$ can be calculated from the formula

$$\mathbf{E}(t, \mathbf{x}) = -\nabla \phi(t, \mathbf{x}) - \frac{\partial}{\partial t} \mathbf{A}(t, \mathbf{x}) \quad (3.11)$$

and $\mathbf{B}(t, \mathbf{x})$ from equation (3.7) above. Hence, it is a matter of taste whether we want to express the laws of electrodynamics in terms of the potentials $\phi(t, \mathbf{x})$ and $\mathbf{A}(t, \mathbf{x})$, or in terms of the fields $\mathbf{E}(t, \mathbf{x})$ and $\mathbf{B}(t, \mathbf{x})$. However, there exists an important difference between the two approaches: in classical electrodynamics the only directly observable quantities are the fields themselves (and quantities derived from them) and not the potentials. On the other hand, the treatment becomes significantly simpler if we use the potentials in our calculations and then, at the final stage, use equation (3.7) and equation (3.11) above to calculate the fields or physical quantities expressed in the fields.

Inserting (3.11) and (3.7) into Maxwell's equations (1.43) on page 14 we obtain, after some simple algebra and the use of equation (1.9) on page 5, the *general*

inhomogeneous wave equations

$$\nabla^2 \phi + \frac{\partial}{\partial t} (\nabla \cdot \mathbf{A}) = -\frac{\rho(t, \mathbf{x})}{\epsilon_0} \quad (3.12a)$$

$$\nabla^2 \mathbf{A} - \frac{1}{c^2} \frac{\partial^2}{\partial t^2} \mathbf{A} - \nabla \left(\nabla \cdot \mathbf{A} + \frac{1}{c^2} \frac{\partial}{\partial t} \phi \right) = -\mu_0 \mathbf{j}(t, \mathbf{x}) \quad (3.12b)$$

These two second order, coupled, partial differential equations, representing in all four scalar equations (one for ϕ and one each for the three components A_1 , A_2 , and A_3 of \mathbf{A}) are completely equivalent to the formulation of electrodynamics in terms of Maxwell's equations, which represent eight scalar first-order, coupled, partial differential equations.

3.3.1 Electromagnetic gauges

Lorentz equations for the electromagnetic potentials

As they stand, Equations (3.12) look complicated and may seem to be of limited use. However, if we write equation (3.7) on the preceding page in the form $\nabla \times \mathbf{A}(t, \mathbf{x}) = \mathbf{B}(t, \mathbf{x})$ we can consider this as a specification of $\nabla \times \mathbf{A}$. But we know from *Helmholtz' theorem* that in order to determine the (spatial behaviour) of \mathbf{A} completely, we must also specify $\nabla \cdot \mathbf{A}$. Since this divergence does not enter the derivation above, *we are free to choose $\nabla \cdot \mathbf{A}$ in whatever way we like and still obtain the same physical results!* This is somewhat analogous to the freedom of adding an arbitrary scalar constant (whose grad vanishes) to the potential energy in classical mechanics and still get the same force.

With a judicious choice of $\nabla \cdot \mathbf{A}$, the calculations can be simplified considerably. Lorentz introduced

$$\nabla \cdot \mathbf{A} + \frac{1}{c^2} \frac{\partial}{\partial t} \phi = 0 \quad (3.13)$$

which is called the *Lorentz gauge condition*, because this choice simplifies the system of *coupled* equations (3.12) above into the following set of *uncoupled* partial differential equations:

$$\square^2 \phi \stackrel{\text{def}}{=} \frac{1}{c^2} \frac{\partial^2}{\partial t^2} \phi - \nabla^2 \phi = \frac{\rho(t, \mathbf{x})}{\epsilon_0} \quad (3.14a)$$

$$\square^2 \mathbf{A} \stackrel{\text{def}}{=} \frac{1}{c^2} \frac{\partial^2}{\partial t^2} \mathbf{A} - \nabla^2 \mathbf{A} = \mu_0 \mathbf{j}(t, \mathbf{x}) \quad (3.14b)$$

where \square^2 is the *d'Alembert operator* discussed in example M.5 on page 185. We shall call (3.14) the *Lorentz equations* for the electromagnetic potentials.

Gauge transformations

We saw in section 3.1 on page 33 and in section 3.2 on page 34 that in electrostatics and magnetostatics we have a certain *mathematical* degree of freedom, up to terms of vanishing gradients and curls, to pick suitable forms for the potentials and still get the same *physical* result. In fact, the way the electromagnetic scalar potential $\phi(t, \mathbf{x})$ and the vector potential $\mathbf{A}(t, \mathbf{x})$ are related to the physically observables gives leeway for similar “manipulation” of them also in electrodynamics. If we transform $\phi(t, \mathbf{x})$ and $\mathbf{A}(t, \mathbf{x})$ *simultaneously* into new ones $\phi'(t, \mathbf{x})$ and $\mathbf{A}'(t, \mathbf{x})$ according to the mapping scheme

$$\phi(t, \mathbf{x}) \rightsquigarrow \phi'(t, \mathbf{x}) = \phi(t, \mathbf{x}) + \frac{\partial}{\partial t} \Gamma(t, \mathbf{x}) \quad (3.15a)$$

$$\mathbf{A}(t, \mathbf{x}) \rightsquigarrow \mathbf{A}'(t, \mathbf{x}) = \mathbf{A}(t, \mathbf{x}) - \nabla \Gamma(t, \mathbf{x}) \quad (3.15b)$$

where $\Gamma(t, \mathbf{x})$ is an arbitrary, differentiable scalar function called the *gauge function*, and insert the transformed potentials into equation (3.11) on page 35 for the electric field and into equation (3.7) on page 35 for the magnetic field, we obtain the transformed fields

$$\begin{aligned} \mathbf{E}' &= -\nabla \phi' - \frac{\partial}{\partial t} \mathbf{A}' = -\nabla \phi - \frac{\partial}{\partial t} \nabla \Gamma - \frac{\partial}{\partial t} \mathbf{A} + \frac{\partial}{\partial t} \nabla \Gamma \\ &= -\nabla \phi - \frac{\partial}{\partial t} \mathbf{A} \end{aligned} \quad (3.16a)$$

$$\mathbf{B}' = \nabla \times \mathbf{A}' = \nabla \times \mathbf{A} - \nabla \times \nabla \Gamma = \nabla \times \mathbf{A} \quad (3.16b)$$

where, once again equation (M.75) on page 188 was used. Comparing these expressions with (3.11) and (3.7) we see that the fields are unaffected by the gauge transformation (3.15). A transformation of the potentials ϕ and \mathbf{A} which leaves the fields, and hence Maxwell’s equations, invariant is called a *gauge transformation*. A physical law which does not change under a gauge transformation is said to be *gauge invariant*. By definition, the fields themselves are, of course, gauge invariant.

The potentials $\phi(t, \mathbf{x})$ and $\mathbf{A}(t, \mathbf{x})$ calculated from (3.12) on the facing page, with an arbitrary choice of $\nabla \cdot \mathbf{A}$, can be further gauge transformed according to (3.15) above. If, in particular, we choose $\nabla \cdot \mathbf{A}$ according to the Lorentz condition, equation (3.13) on the facing page, and apply the gauge transformation (3.15) on the resulting Lorentz equations (3.14) on the preceding page, these equations will be transformed into

$$\frac{1}{c^2} \frac{\partial^2}{\partial t^2} \phi - \nabla^2 \phi + \frac{\partial}{\partial t} \left(\frac{1}{c^2} \frac{\partial^2}{\partial t^2} \Gamma - \nabla^2 \Gamma \right) = \frac{\rho(t, \mathbf{x})}{\epsilon_0} \quad (3.17a)$$

$$\frac{1}{c^2} \frac{\partial^2}{\partial t^2} \mathbf{A} - \nabla^2 \mathbf{A} - \nabla \left(\frac{1}{c^2} \frac{\partial^2}{\partial t^2} \Gamma - \nabla^2 \Gamma \right) = \mu_0 \mathbf{j}(t, \mathbf{x}) \quad (3.17b)$$

We notice that if we require that the gauge function $\Gamma(t, \mathbf{x})$ itself be restricted to fulfil the wave equation

$$\frac{1}{c^2} \frac{\partial^2}{\partial t^2} \Gamma - \nabla^2 \Gamma = 0 \quad (3.18)$$

these transformed Lorentz equations will keep their original form. The set of potentials which have been gauge transformed according to equation (3.15) on the preceding page with a gauge function $\Gamma(t, \mathbf{x})$ which is restricted to fulfil equation (3.18) above, *i.e.*, those gauge transformed potentials for which the Lorentz equations (3.14) are invariant, comprises the *Lorentz gauge*.

Other useful gauges are

- The *radiation gauge*, also known as the *transverse gauge*, defined by $\nabla \cdot \mathbf{A} = 0$.
- The *Coulomb gauge*, defined by $\phi = 0, \nabla \cdot \mathbf{A} = 0$.
- The *temporal gauge*, also known as the *Hamilton gauge*, defined by $\phi = 0$.
- The *axial gauge*, defined by $A_3 = 0$.

The process of choosing a particular gauge condition is referred to as *gauge fixing*.

3.3.2 Solution of the Lorentz equations for the electromagnetic potentials

As we see, the Lorentz equations (3.14) on page 36 for $\phi(t, \mathbf{x})$ and $\mathbf{A}(t, \mathbf{x})$ represent a set of uncoupled equations involving four scalar equations (one equation for ϕ and one equation for each of the three components of \mathbf{A}). Each of these four scalar equations is an *inhomogeneous wave equation* of the following generic form:

$$\square^2 \Psi(t, \mathbf{x}) = f(t, \mathbf{x}) \quad (3.19)$$

where Ψ is a shorthand for either ϕ or one of the vector components of \mathbf{A} , and f is the pertinent generic source component.

We assume that our sources are well-behaved enough in time t so that the *Fourier*

transform pair for the generic source function

$$\mathcal{F}^{-1}[f_\omega(\mathbf{x})] \stackrel{\text{def}}{=} f(t, \mathbf{x}) = \int_{-\infty}^{\infty} f_\omega(\mathbf{x}) e^{-i\omega t} d\omega \quad (3.20a)$$

$$\mathcal{F}[f(t, \mathbf{x})] \stackrel{\text{def}}{=} f_\omega(\mathbf{x}) = \frac{1}{2\pi} \int_{-\infty}^{\infty} f(t, \mathbf{x}) e^{i\omega t} dt \quad (3.20b)$$

exists, and that the same is true for the generic potential component:

$$\Psi(t, \mathbf{x}) = \int_{-\infty}^{\infty} \Psi_\omega(\mathbf{x}) e^{-i\omega t} d\omega \quad (3.21a)$$

$$\Psi_\omega(\mathbf{x}) = \frac{1}{2\pi} \int_{-\infty}^{\infty} \Psi(t, \mathbf{x}) e^{i\omega t} dt \quad (3.21b)$$

Inserting the Fourier representations (3.20a) and (3.21a) into equation (3.19) on the preceding page, and using the vacuum dispersion relation for electromagnetic waves

$$\omega = ck \quad (3.22)$$

the generic 3D inhomogeneous wave equation equation (3.19) on the facing page turns into

$$\nabla^2 \Psi_\omega(\mathbf{x}) + k^2 \Psi_\omega(\mathbf{x}) = -f_\omega(\mathbf{x}) \quad (3.23)$$

which is a 3D *inhomogeneous time-independent wave equation*, often called the 3D *inhomogeneous Helmholtz equation*.

As postulated by *Huygen's principle*, each point on a wave front acts as a point source for spherical waves which form a new wave from a superposition of the individual waves from each of the point sources on the old wave front. The solution of (3.23) can therefore be expressed as a superposition of solutions of an equation where the source term has been replaced by a point source:

$$\nabla^2 G(\mathbf{x}, \mathbf{x}') + k^2 G(\mathbf{x}, \mathbf{x}') = -\delta(\mathbf{x} - \mathbf{x}') \quad (3.24)$$

and the solution of equation (3.23) above which corresponds to the frequency ω is given by the superposition

$$\Psi_\omega = \int f_\omega(\mathbf{x}') G(\mathbf{x}, \mathbf{x}') d^3x' \quad (3.25)$$

The function $G(\mathbf{x}, \mathbf{x}')$ is called the *Green's function* or the *propagator*.

In equation (3.24), the Dirac generalised function $\delta(\mathbf{x} - \mathbf{x}')$, which represents the point source, depends only on $\mathbf{x} - \mathbf{x}'$ and there is no angular dependence in the equation. Hence, the solution can only be dependent on $r = |\mathbf{x} - \mathbf{x}'|$. If we interpret r as the radial coordinate in a spherically polar coordinate system, the

“spherically symmetric” $G(r)$ is given by the solution of

$$\frac{d^2}{dr^2}(rG) + k^2(rG) = -r\delta(r) \quad (3.26)$$

Away from $r = |\mathbf{x} - \mathbf{x}'| = 0$, *i.e.*, away from the source point \mathbf{x}' , this equation takes the form

$$\frac{d^2}{dr^2}(rG) + k^2(rG) = 0 \quad (3.27)$$

with the well-known general solution

$$G = C^+ \frac{e^{ikr}}{r} + C^- \frac{e^{-ikr}}{r} \equiv C^+ \frac{e^{ik|\mathbf{x}-\mathbf{x}'|}}{|\mathbf{x}-\mathbf{x}'|} + C^- \frac{e^{-ik|\mathbf{x}-\mathbf{x}'|}}{|\mathbf{x}-\mathbf{x}'|} \quad (3.28)$$

where C^\pm are constants.

In order to evaluate the constants C^\pm , we insert the general solution, equation (3.28) above, into equation (3.24) on the preceding page and integrate over a small volume around $r = |\mathbf{x} - \mathbf{x}'| = 0$. Since

$$G(|\mathbf{x} - \mathbf{x}'|) \sim C^+ \frac{1}{|\mathbf{x} - \mathbf{x}'|} + C^- \frac{1}{|\mathbf{x} - \mathbf{x}'|}, \quad |\mathbf{x} - \mathbf{x}'| \rightarrow 0 \quad (3.29)$$

equation (3.24) on the previous page can under this assumption be approximated by

$$\begin{aligned} (C^+ + C^-) \int \nabla^2 \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right) d^3x' + k^2 (C^+ + C^-) \int \frac{1}{|\mathbf{x} - \mathbf{x}'|} d^3x' \\ = - \int \delta(|\mathbf{x} - \mathbf{x}'|) d^3x' \end{aligned} \quad (3.30)$$

In virtue of the fact that the volume element d^3x' in spherical polar coordinates is proportional to $|\mathbf{x} - \mathbf{x}'|^2$, the second integral vanishes when $|\mathbf{x} - \mathbf{x}'| \rightarrow 0$. Furthermore, from equation (M.70) on page 187, we find that the integrand in the first integral can be written as $-4\pi\delta(|\mathbf{x} - \mathbf{x}'|)$ and, hence, that

$$C^+ + C^- = \frac{1}{4\pi} \quad (3.31)$$

Insertion of the general solution equation (3.28) into equation (3.25) on the previous page gives

$$\Psi_\omega(\mathbf{x}) = C^+ \int f_\omega(\mathbf{x}') \frac{e^{ik|\mathbf{x}-\mathbf{x}'|}}{|\mathbf{x}-\mathbf{x}'|} d^3x' + C^- \int f_\omega(\mathbf{x}') \frac{e^{-ik|\mathbf{x}-\mathbf{x}'|}}{|\mathbf{x}-\mathbf{x}'|} d^3x' \quad (3.32)$$

The Fourier transform to ordinary t domain of this is obtained by inserting the

above expression for $\Psi_\omega(\mathbf{x})$ into equation (3.21a) on page 39:

$$\begin{aligned}\Psi(t, \mathbf{x}) &= C^+ \iint f_\omega(\mathbf{x}') \frac{\exp\left[-i\omega\left(t - \frac{k|\mathbf{x}-\mathbf{x}'|}{\omega}\right)\right]}{|\mathbf{x}-\mathbf{x}'|} d\omega d^3x' \\ &+ C^- \iint f_\omega(\mathbf{x}') \frac{\exp\left[-i\omega\left(t + \frac{k|\mathbf{x}-\mathbf{x}'|}{\omega}\right)\right]}{|\mathbf{x}-\mathbf{x}'|} d\omega d^3x'\end{aligned}\quad (3.33)$$

If we introduce the *retarded time* t'_{ret} and the *advanced time* t'_{adv} in the following way [using the fact that in vacuum $k/\omega = 1/c$, according to equation (3.22) on page 39]:

$$t'_{\text{ret}} = t'_{\text{ret}}(t, |\mathbf{x}-\mathbf{x}'|) = t - \frac{k|\mathbf{x}-\mathbf{x}'|}{\omega} = t - \frac{|\mathbf{x}-\mathbf{x}'|}{c} \quad (3.34a)$$

$$t'_{\text{adv}} = t'_{\text{adv}}(t, |\mathbf{x}-\mathbf{x}'|) = t + \frac{k|\mathbf{x}-\mathbf{x}'|}{\omega} = t + \frac{|\mathbf{x}-\mathbf{x}'|}{c} \quad (3.34b)$$

and use equation (3.20a) on page 39, we obtain

$$\Psi(t, \mathbf{x}) = C^+ \int \frac{f(t'_{\text{ret}}, \mathbf{x}')}{|\mathbf{x}-\mathbf{x}'|} d^3x' + C^- \int \frac{f(t'_{\text{adv}}, \mathbf{x}')}{|\mathbf{x}-\mathbf{x}'|} d^3x' \quad (3.35)$$

This is a solution to the generic inhomogeneous wave equation for the potential components equation (3.19) on page 38. We note that the solution at time t at the field point \mathbf{x} is dependent on the behaviour at other times t' of the source at \mathbf{x}' and that both retarded and advanced t' are mathematically acceptable solutions. However, if we assume that causality requires that the potential at (t, \mathbf{x}) is set up by the source at an earlier time, *i.e.*, at $(t'_{\text{ret}}, \mathbf{x}')$, we must in equation (3.35) set $C^- = 0$ and therefore, according to equation (3.31) on the preceding page, $C^+ = 1/(4\pi)$.

The retarded potentials

From the above discussion on the solution of the inhomogeneous wave equation we conclude that under the assumption of causality the electromagnetic potentials in vacuum can be written

$$\phi(t, \mathbf{x}) = \frac{1}{4\pi\epsilon_0} \int \frac{\rho(t'_{\text{ret}}, \mathbf{x}')}{|\mathbf{x}-\mathbf{x}'|} d^3x' \quad (3.36a)$$

$$\mathbf{A}(t, \mathbf{x}) = \frac{\mu_0}{4\pi} \int \frac{\mathbf{j}(t'_{\text{ret}}, \mathbf{x}')}{|\mathbf{x}-\mathbf{x}'|} d^3x' \quad (3.36b)$$

Since these *retarded potentials* were obtained as solutions to the Lorentz equations (3.14) on page 36 they are valid in the Lorentz gauge but may be gauge transformed according to the scheme described in subsection 3.3.1 on page 37. As they stand, we shall use them frequently in the following.

- [1] L. D. Fadeev and A. A. Slavnov. *Gauge Fields: Introduction to Quantum Theory*. Number 50 in Frontiers in Physics: A Lecture Note and Reprint Series. Benjamin/Cummings Publishing Company, Inc., Reading, MA ..., 1980. ISBN 0-8053-9016-2.
- [2] Mike Guidry. *Gauge Field Theories: An Introduction with Applications*. Wiley & Sons, Inc., New York, NY ..., 1991. ISBN 0-471-63117-5.
- [3] John D. Jackson. *Classical Electrodynamics*. Wiley & Sons, Inc., New York, NY ..., third edition, 1999. ISBN 0-471-30932-X.
- [4] Wolfgang K. H. Panofsky and Melba Phillips. *Classical Electricity and Magnetism*. Addison-Wesley Publishing Company, Inc., Reading, MA ..., third edition, 1962. ISBN 0-201-05702-6.

The Electromagnetic Fields

While, in principle, the electric and magnetic fields can be calculated from the Maxwell equations in chapter 1, or even from the wave equations in chapter 2, it is often physically more lucid to calculate them from the electromagnetic potentials derived in chapter 3. In this Chapter we will derive the electric and magnetic fields from the potentials.

We recall that in order to find the solution (3.35) for the generic inhomogeneous wave equation (3.19) on page 38 we presupposed the existence of a Fourier transform pair (3.20a) on page 39 for the generic source term

$$f(t, \mathbf{x}) = \int_{-\infty}^{\infty} f_{\omega}(\mathbf{x}) e^{-i\omega t} d\omega \quad (4.1a)$$

$$f_{\omega}(\mathbf{x}) = \frac{1}{2\pi} \int_{-\infty}^{\infty} f(t, \mathbf{x}) e^{i\omega t} dt \quad (4.1b)$$

and for the generic potential component

$$\Psi(t, \mathbf{x}) = \int_{-\infty}^{\infty} \Psi_{\omega}(\mathbf{x}) e^{-i\omega t} d\omega \quad (4.2a)$$

$$\Psi_{\omega}(\mathbf{x}) = \frac{1}{2\pi} \int_{-\infty}^{\infty} \Psi(t, \mathbf{x}) e^{i\omega t} dt \quad (4.2b)$$

That such transform pairs exists is true for most physical variables which are not strictly monotonically increasing and decreasing with time. For charges and currents varying in time we can therefore, without loss of generality, work with individual Fourier components. Strictly speaking, the existence of a single Fourier component assumes a *monochromatic* source (*i.e.*, a source containing only one single frequency component), which in turn requires that the electric and magnetic fields exist for infinitely long times. However, by taking the proper limits, we can still use this approach even for sources and fields of finite duration.

This is the method we shall utilise in this Chapter in order to derive the electric and magnetic fields in vacuum from arbitrary given charge densities $\rho(t, \mathbf{x})$ and

current densities $\mathbf{j}(t, \mathbf{x})$, defined by the Fourier transform pairs

$$\rho(t, \mathbf{x}) = \int_{-\infty}^{\infty} \rho_{\omega}(\mathbf{x}) e^{-i\omega t} d\omega \quad (4.3a)$$

$$\rho_{\omega}(\mathbf{x}) = \frac{1}{2\pi} \int_{-\infty}^{\infty} \rho(t, \mathbf{x}) e^{i\omega t} dt \quad (4.3b)$$

and

$$\mathbf{j}(t, \mathbf{x}) = \int_{-\infty}^{\infty} \mathbf{j}_{\omega}(\mathbf{x}) e^{-i\omega t} d\omega \quad (4.4a)$$

$$\mathbf{j}_{\omega}(\mathbf{x}) = \frac{1}{2\pi} \int_{-\infty}^{\infty} \mathbf{j}(t, \mathbf{x}) e^{i\omega t} dt \quad (4.4b)$$

under the assumption that only *retarded* potentials produce physically acceptable solutions.¹

The Fourier transform pair for the retarded vector potential can then be written

$$\phi(t, \mathbf{x}) = \int_{-\infty}^{\infty} \phi_{\omega}(\mathbf{x}) e^{-i\omega t} d\omega \quad (4.5a)$$

$$\phi_{\omega}(\mathbf{x}) = \frac{1}{2\pi} \int_{-\infty}^{\infty} \phi(t, \mathbf{x}) e^{i\omega t} dt = \frac{1}{4\pi\epsilon_0} \int \rho_{\omega}(\mathbf{x}') \frac{e^{ik|\mathbf{x}-\mathbf{x}'|}}{|\mathbf{x}-\mathbf{x}'|} d^3x' \quad (4.5b)$$

where in the last step, we made use of the explicit expression for the Fourier transform of the generic potential component $\Psi_{\omega}(\mathbf{x})$, equation (3.32) on page 40. Similarly, the following Fourier transform pair for the vector potential must exist:

$$\mathbf{A}(t, \mathbf{x}) = \int_{-\infty}^{\infty} \mathbf{A}_{\omega}(\mathbf{x}) e^{-i\omega t} d\omega \quad (4.6a)$$

$$\mathbf{A}_{\omega}(\mathbf{x}) = \frac{1}{2\pi} \int_{-\infty}^{\infty} \mathbf{A}(t, \mathbf{x}) e^{i\omega t} dt = \frac{\mu_0}{4\pi} \int \mathbf{j}_{\omega}(\mathbf{x}') \frac{e^{ik|\mathbf{x}-\mathbf{x}'|}}{|\mathbf{x}-\mathbf{x}'|} d^3x' \quad (4.6b)$$

Clearly, we must require that

$$\mathbf{A}_{\omega} = \mathbf{A}_{-\omega}^*, \quad \phi_{\omega} = \phi_{-\omega}^* \quad (4.7)$$

in order that all physical quantities be real. Similar transform pairs and requirements of real valuedness exist for the fields themselves.

In the limit that the sources can be considered monochromatic containing only

¹In fact, John A. Wheeler and Richard P. Feynman derived in 1945 a fully self-consistent electrody-
namics using both the retarded and the advanced potentials [?]; See also [?].

one single frequency ω_0 , we have the much simpler expressions

$$\rho(t, \mathbf{x}) = \rho_0(\mathbf{x})e^{-i\omega_0 t} \quad (4.8a)$$

$$\mathbf{j}(t, \mathbf{x}) = \mathbf{j}_0(\mathbf{x})e^{-i\omega_0 t} \quad (4.8b)$$

$$\phi(t, \mathbf{x}) = \phi_0(\mathbf{x})e^{-i\omega_0 t} \quad (4.8c)$$

$$\mathbf{A}(t, \mathbf{x}) = \mathbf{A}_0(\mathbf{x})e^{-i\omega_0 t} \quad (4.8d)$$

where again the real-valuedness of all these quantities is implied. As discussed above, we can safely assume that all formulae derived for a general Fourier representation of the source (general distribution of frequencies in the source) are valid for these simple limiting cases. We note that in this context, we can make the formal identification $\rho_\omega = \rho_0\delta(\omega - \omega_0)$, $\mathbf{j}_\omega = \mathbf{j}_0\delta(\omega - \omega_0)$ *etc.*, and that we therefore, without any loss of stringence, let ρ_0 mean the same as the Fourier amplitude ρ_ω and so on.

4.1 The magnetic field

Let us now compute the magnetic field from the vector potential, defined by Equation (4.6a) and equation (4.6b) on the preceding page, and formula (3.7) on page 35:

$$\mathbf{B}(t, \mathbf{x}) = \nabla \times \mathbf{A}(t, \mathbf{x}) \quad (4.9)$$

The calculations are much simplified if we work in ω space and, at the final stage, Fourier transform back to ordinary t space. We are working in the Lorentz gauge and note that in ω space the Lorentz condition, equation (3.13) on page 36, takes the form

$$\nabla \cdot \mathbf{A}_\omega - i\frac{k}{c}\phi_\omega = 0 \quad (4.10)$$

which provides a relation between (the Fourier transforms of) the vector and scalar potentials.

Using the Fourier transformed version of equation (4.9) and equation (4.6b) on

page 46, we obtain

$$\mathbf{B}_\omega(\mathbf{x}) = \nabla \times \mathbf{A}_\omega(\mathbf{x}) = \frac{\mu_0}{4\pi} \nabla \times \int_V \mathbf{j}_\omega(\mathbf{x}') \frac{e^{ik|\mathbf{x}-\mathbf{x}'|}}{|\mathbf{x}-\mathbf{x}'|} d^3x' \quad (4.11)$$

Using formula (F.62) on page 167, we can rewrite this as

$$\begin{aligned} \mathbf{B}_\omega(\mathbf{x}) &= -\frac{\mu_0}{4\pi} \int_V \mathbf{j}_\omega(\mathbf{x}') \times \left[\nabla \left(\frac{e^{ik|\mathbf{x}-\mathbf{x}'|}}{|\mathbf{x}-\mathbf{x}'|} \right) \right] d^3x' \\ &= -\frac{\mu_0}{4\pi} \left[\int_V \mathbf{j}_\omega(\mathbf{x}') \times \left(-\frac{\mathbf{x}-\mathbf{x}'}{|\mathbf{x}-\mathbf{x}'|^3} \right) e^{ik|\mathbf{x}-\mathbf{x}'|} d^3x' \right. \\ &\quad \left. + \int_V \mathbf{j}_\omega(\mathbf{x}') \times \left(ik \frac{\mathbf{x}-\mathbf{x}'}{|\mathbf{x}-\mathbf{x}'|} e^{ik|\mathbf{x}-\mathbf{x}'|} \right) \frac{1}{|\mathbf{x}-\mathbf{x}'|} d^3x' \right] \\ &= \frac{\mu_0}{4\pi} \left[\int_V \frac{\mathbf{j}_\omega(\mathbf{x}') e^{ik|\mathbf{x}-\mathbf{x}'|} \times (\mathbf{x}-\mathbf{x}')}{|\mathbf{x}-\mathbf{x}'|^3} d^3x' \right. \\ &\quad \left. + \int_V \frac{(-ik) \mathbf{j}_\omega(\mathbf{x}') e^{ik|\mathbf{x}-\mathbf{x}'|} \times (\mathbf{x}-\mathbf{x}')}{|\mathbf{x}-\mathbf{x}'|^2} d^3x' \right] \end{aligned} \quad (4.12)$$

From this expression for the magnetic field in the frequency (ω) domain, we obtain the magnetic field in the temporal (t) domain by taking the inverse Fourier transform (using the identity $-ik = -i\omega/c$):

$$\begin{aligned} \mathbf{B}(t, \mathbf{x}) &= \int_{-\infty}^{\infty} \mathbf{B}_\omega(\mathbf{x}) e^{-i\omega t} d\omega \\ &= \frac{\mu_0}{4\pi} \left\{ \int_V \frac{\left[\int \mathbf{j}_\omega(\mathbf{x}') e^{i(k|\mathbf{x}-\mathbf{x}'|-\omega t)} d\omega \right] \times (\mathbf{x}-\mathbf{x}')}{|\mathbf{x}-\mathbf{x}'|^3} d^3x' \right. \\ &\quad \left. + \frac{1}{c} \int_V \frac{\left[\int (-i\omega) \mathbf{j}_\omega(\mathbf{x}') e^{i(k|\mathbf{x}-\mathbf{x}'|-\omega t)} d\omega \right] \times (\mathbf{x}-\mathbf{x}')}{|\mathbf{x}-\mathbf{x}'|^2} d^3x' \right\} \\ &= \frac{\mu_0}{4\pi} \underbrace{\int_V \frac{\mathbf{j}(t'_{\text{ret}}, \mathbf{x}') \times (\mathbf{x}-\mathbf{x}')}{|\mathbf{x}-\mathbf{x}'|^3} d^3x'}_{\text{Induction field}} \\ &\quad + \frac{\mu_0}{4\pi c} \underbrace{\int_V \frac{\dot{\mathbf{j}}(t'_{\text{ret}}, \mathbf{x}') \times (\mathbf{x}-\mathbf{x}')}{|\mathbf{x}-\mathbf{x}'|^2} d^3x'}_{\text{Radiation field}} \end{aligned} \quad (4.13)$$

where

$$\mathbf{j}(t'_{\text{ret}}, \mathbf{x}') \stackrel{\text{def}}{=} \left(\frac{\partial \mathbf{j}}{\partial t} \right)_{t=t'_{\text{ret}}} \quad (4.14)$$

The first term, the *induction field*, dominates near the current source but falls off rapidly with distance from it, while the second term, the *radiation field* or the *far field*, dominates at large distances and represents energy that is transported out to infinity. Note how the spatial derivatives (∇) gave rise to a time derivative ($\dot{}$)!

4.2 The electric field

To calculate the electric field, we use the Fourier transformed version of formula (3.11) on page 35, inserting Equations (4.5b) and (4.6b) as the explicit expressions for the Fourier transforms of ϕ and \mathbf{A} :

$$\begin{aligned} \mathbf{E}_\omega(\mathbf{x}) &= -\nabla\phi_\omega(\mathbf{x}) + i\omega\mathbf{A}_\omega(\mathbf{x}) \\ &= -\frac{1}{4\pi\epsilon_0}\nabla\int_V\rho_\omega(\mathbf{x}')\frac{e^{ik|\mathbf{x}-\mathbf{x}'|}}{|\mathbf{x}-\mathbf{x}'|}d^3x' + \frac{i\mu_0\omega}{4\pi}\int_V\mathbf{j}_\omega(\mathbf{x}')\frac{e^{ik|\mathbf{x}-\mathbf{x}'|}}{|\mathbf{x}-\mathbf{x}'|}d^3x' \\ &= \frac{1}{4\pi\epsilon_0}\left[\int_V\frac{\rho_\omega(\mathbf{x}')e^{ik|\mathbf{x}-\mathbf{x}'|}(\mathbf{x}-\mathbf{x}')}{|\mathbf{x}-\mathbf{x}'|^3}d^3x' \right. \\ &\quad \left. - ik\int_V\left(\frac{\rho_\omega(\mathbf{x}')(\mathbf{x}-\mathbf{x}')}{|\mathbf{x}-\mathbf{x}'|} - \frac{\mathbf{j}_\omega(\mathbf{x}')}{c}\right)\frac{e^{ik|\mathbf{x}-\mathbf{x}'|}}{|\mathbf{x}-\mathbf{x}'|}d^3x'\right] \end{aligned} \quad (4.15)$$

Using the Fourier transform of the continuity equation (1.21) on page 9

$$\nabla' \cdot \mathbf{j}_\omega(\mathbf{x}') - i\omega\rho_\omega(\mathbf{x}') = 0 \quad (4.16)$$

we see that we can express ρ_ω in terms of \mathbf{j}_ω as follows

$$\rho_\omega(\mathbf{x}') = -\frac{i}{\omega}\nabla' \cdot \mathbf{j}_\omega(\mathbf{x}') \quad (4.17)$$

Doing so in the last term of equation (4.15) above, and also using the fact that $k = \omega/c$, we can rewrite this Equation as

$$\begin{aligned} \mathbf{E}_\omega(\mathbf{x}) &= \frac{1}{4\pi\epsilon_0}\left[\int_V\frac{\rho_\omega(\mathbf{x}')e^{ik|\mathbf{x}-\mathbf{x}'|}(\mathbf{x}-\mathbf{x}')}{|\mathbf{x}-\mathbf{x}'|^3}d^3x' \right. \\ &\quad \left. - \frac{1}{c}\int_V\left(\frac{[\nabla' \cdot \mathbf{j}_\omega(\mathbf{x}')](\mathbf{x}-\mathbf{x}')}{|\mathbf{x}-\mathbf{x}'|} - ik\mathbf{j}_\omega(\mathbf{x}')\right)\frac{e^{ik|\mathbf{x}-\mathbf{x}'|}}{|\mathbf{x}-\mathbf{x}'|}d^3x'\right] \\ &\quad \underbrace{\hspace{10em}}_{\mathbf{I}_\omega} \end{aligned} \quad (4.18)$$

The last integral can be further rewritten in the following way:

$$\begin{aligned}
\mathbf{I}_\omega &= \int_V \left(\frac{[\nabla' \cdot \mathbf{j}_\omega(\mathbf{x}')](\mathbf{x} - \mathbf{x}')}{|\mathbf{x} - \mathbf{x}'|} - ik\mathbf{j}_\omega(\mathbf{x}') \right) \frac{e^{ik|\mathbf{x}-\mathbf{x}'|}}{|\mathbf{x} - \mathbf{x}'|} d^3x' \\
&= \int_V \left(\frac{\partial j_{\omega m}}{\partial x'_m} \frac{x_l - x'_l}{|\mathbf{x} - \mathbf{x}'|} - ikj_{\omega l}(\mathbf{x}') \right) \hat{\mathbf{x}}_l \frac{e^{ik|\mathbf{x}-\mathbf{x}'|}}{|\mathbf{x} - \mathbf{x}'|} d^3x' \quad (4.19)
\end{aligned}$$

But, since

$$\begin{aligned}
\frac{\partial}{\partial x'_m} \left(j_{\omega m} \frac{x_l - x'_l}{|\mathbf{x} - \mathbf{x}'|^2} e^{ik|\mathbf{x}-\mathbf{x}'|} \right) &= \left(\frac{\partial j_{\omega m}}{\partial x'_m} \right) \frac{x_l - x'_l}{|\mathbf{x} - \mathbf{x}'|^2} e^{ik|\mathbf{x}-\mathbf{x}'|} \\
&\quad + j_{\omega m} \frac{\partial}{\partial x'_m} \left(\frac{x_l - x'_l}{|\mathbf{x} - \mathbf{x}'|^2} e^{ik|\mathbf{x}-\mathbf{x}'|} \right) \quad (4.20)
\end{aligned}$$

we can rewrite \mathbf{I}_ω as

$$\begin{aligned}
\mathbf{I}_\omega &= - \int_V \left[j_{\omega m} \frac{\partial}{\partial x'_m} \left(\frac{x_l - x'_l}{|\mathbf{x} - \mathbf{x}'|^2} \hat{\mathbf{x}}_l e^{ik|\mathbf{x}-\mathbf{x}'|} \right) + ik\mathbf{j}_\omega \frac{e^{ik|\mathbf{x}-\mathbf{x}'|}}{|\mathbf{x} - \mathbf{x}'|} \right] d^3x' \\
&\quad + \int_V \frac{\partial}{\partial x'_m} \left(j_{\omega m} \frac{x_l - x'_l}{|\mathbf{x} - \mathbf{x}'|^2} \hat{\mathbf{x}}_l e^{ik|\mathbf{x}-\mathbf{x}'|} \right) d^3x' \quad (4.21)
\end{aligned}$$

where, according to Gauss's theorem, the last term vanishes if \mathbf{j}_ω is assumed to be limited and tends to zero at large distances. Further evaluation of the derivative in the first term makes it possible to write

$$\begin{aligned}
\mathbf{I}_\omega &= - \int_V \left(-\mathbf{j}_\omega \frac{e^{ik|\mathbf{x}-\mathbf{x}'|}}{|\mathbf{x} - \mathbf{x}'|^2} + \frac{2}{|\mathbf{x} - \mathbf{x}'|^4} [\mathbf{j}_\omega \cdot (\mathbf{x} - \mathbf{x}')] (\mathbf{x} - \mathbf{x}') e^{ik|\mathbf{x}-\mathbf{x}'|} \right) d^3x' \\
&\quad - ik \int_V \left(-\frac{[\mathbf{j}_\omega \cdot (\mathbf{x} - \mathbf{x}')] (\mathbf{x} - \mathbf{x}')}{|\mathbf{x} - \mathbf{x}'|^3} e^{ik|\mathbf{x}-\mathbf{x}'|} + \mathbf{j}_\omega \frac{e^{ik|\mathbf{x}-\mathbf{x}'|}}{|\mathbf{x} - \mathbf{x}'|} \right) d^3x' \quad (4.22)
\end{aligned}$$

Using the triple product ‘‘bac-cab’’ formula (F.56) on page 167 backwards, and inserting the resulting expression for \mathbf{I}_ω into equation (4.18) on the previous page, we arrive at the following final expression for the Fourier transform of the total \mathbf{E} -field:

$$\begin{aligned}
\mathbf{E}_\omega(\mathbf{x}) &= -\frac{1}{4\pi\epsilon_0} \nabla \int_V \rho_\omega(\mathbf{x}') \frac{e^{ik|\mathbf{x}-\mathbf{x}'|}}{|\mathbf{x}-\mathbf{x}'|} d^3x' + \frac{i\mu_0\omega}{4\pi} \int_V \mathbf{j}_\omega(\mathbf{x}') \frac{e^{ik|\mathbf{x}-\mathbf{x}'|}}{|\mathbf{x}-\mathbf{x}'|} d^3x' \\
&= \frac{1}{4\pi\epsilon_0} \left[\int_V \frac{\rho_\omega(\mathbf{x}') e^{ik|\mathbf{x}-\mathbf{x}'|} (\mathbf{x}-\mathbf{x}')}{|\mathbf{x}-\mathbf{x}'|^3} d^3x' \right. \\
&\quad + \frac{1}{c} \int_V \frac{[\mathbf{j}_\omega(\mathbf{x}') e^{ik|\mathbf{x}-\mathbf{x}'|} \cdot (\mathbf{x}-\mathbf{x}')] (\mathbf{x}-\mathbf{x}')}{|\mathbf{x}-\mathbf{x}'|^4} d^3x' \\
&\quad + \frac{1}{c} \int_V \frac{[\mathbf{j}_\omega(\mathbf{x}') e^{ik|\mathbf{x}-\mathbf{x}'|} \times (\mathbf{x}-\mathbf{x}')] \times (\mathbf{x}-\mathbf{x}')}{|\mathbf{x}-\mathbf{x}'|^4} d^3x' \\
&\quad \left. - \frac{ik}{c} \int_V \frac{[\mathbf{j}_\omega(\mathbf{x}') e^{ik|\mathbf{x}-\mathbf{x}'|} \times (\mathbf{x}-\mathbf{x}')] \times (\mathbf{x}-\mathbf{x}')}{|\mathbf{x}-\mathbf{x}'|^3} d^3x' \right] \quad (4.23)
\end{aligned}$$

Taking the inverse Fourier transform of equation (4.23), once again using the vacuum relation $\omega = kc$, we find, at last, the expression in time domain for the total electric field:

$$\begin{aligned}
\mathbf{E}(t, \mathbf{x}) &= \int_{-\infty}^{\infty} \mathbf{E}_\omega(\mathbf{x}) e^{-i\omega t} d\omega \\
&= \underbrace{\frac{1}{4\pi\epsilon_0} \int_V \frac{\rho(t'_{\text{ret}}, \mathbf{x}') (\mathbf{x}-\mathbf{x}')}{|\mathbf{x}-\mathbf{x}'|^3} d^3x'}_{\text{Retarded Coulomb field}} \\
&\quad + \underbrace{\frac{1}{4\pi\epsilon_0 c} \int_V \frac{[\mathbf{j}(t'_{\text{ret}}, \mathbf{x}') \cdot (\mathbf{x}-\mathbf{x}')] (\mathbf{x}-\mathbf{x}')}{|\mathbf{x}-\mathbf{x}'|^4} d^3x'}_{\text{Intermediate field}} \\
&\quad + \underbrace{\frac{1}{4\pi\epsilon_0 c} \int_V \frac{[\mathbf{j}(t'_{\text{ret}}, \mathbf{x}') \times (\mathbf{x}-\mathbf{x}')] \times (\mathbf{x}-\mathbf{x}')}{|\mathbf{x}-\mathbf{x}'|^4} d^3x'}_{\text{Intermediate field}} \\
&\quad + \underbrace{\frac{1}{4\pi\epsilon_0 c^2} \int_V \frac{[\dot{\mathbf{j}}(t'_{\text{ret}}, \mathbf{x}') \times (\mathbf{x}-\mathbf{x}')] \times (\mathbf{x}-\mathbf{x}')}{|\mathbf{x}-\mathbf{x}'|^3} d^3x'}_{\text{Radiation field}} \quad (4.24)
\end{aligned}$$

Here, the first term represents the *retarded Coulomb field* and the last term represents the *radiation field* which carries energy over very large distances. The other two terms represent an *intermediate field* which contributes only in the *near zone* and must be taken into account there.

With this we have achieved our goal of finding closed-form analytic expressions for the electric and magnetic fields when the sources of the fields are completely

arbitrary, prescribed distributions of charges and currents. The only assumption made is that the advanced potentials have been discarded; recall the discussion following equation (3.35) on page 41 in chapter 3.

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Relativistic Electrodynamics

We saw in chapter 3 how the derivation of the electrodynamic potentials led, in a most natural way, to the introduction of a characteristic, finite speed of propagation that equals the speed of light $c = 1/\sqrt{\epsilon_0\mu_0}$ and which can be considered as a constant of nature. To take this finite speed of propagation of information into account, and to ensure that our laws of physics be independent of any specific coordinate frame, requires a treatment of electrodynamics in a relativistically covariant (coordinate independent) form. This is the object of the current chapter.

5.1 The special theory of relativity

An *inertial system*, or *inertial reference frame*, is a system of reference, or rigid coordinate system, in which the *law of inertia* (*Galileo's law*, *Newton's first law*) holds. In other words, an inertial system is a system in which free bodies move uniformly and do not experience any acceleration. The *special theory of relativity* describes how physical processes are interrelated when observed in different inertial systems in uniform, rectilinear motion relative to each other and is based on two postulates:

Postulate 1 (*Relativity principle*; Poincaré, 1905) *All laws of physics (except the laws of gravitation) are independent of the uniform translational motion of the system on which they operate.*

Postulate 2 (Einstein, 1905) *The velocity of light is independent of the motion of the source.*

A consequence of the first postulate is that all geometrical objects (vectors, tensors) in an equation describing a physical process must transform in a *covariant* manner, *i.e.*, in the same way.

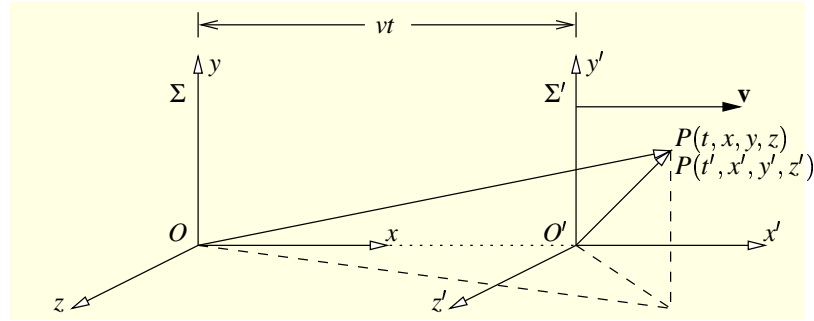


Figure 5.1. Two inertial systems Σ and Σ' in relative motion with velocity \mathbf{v} along the $x = x'$ axis. At time $t = t' = 0$ the origin O' of Σ' coincided with the origin O of Σ . At time t , the inertial system Σ' has been translated a distance vt along the x axis in Σ . An event represented by $P(t, x, y, z)$ in Σ is represented by $P(t', x', y', z')$ in Σ' .

5.1.1 The Lorentz transformation

Let us consider two three-dimensional inertial systems Σ and Σ' in vacuum which are in rectilinear motion relative to each other in such a way that Σ' moves with constant velocity \mathbf{v} along the x axis of the Σ system. The times and the spatial coordinates as measured in the two systems are t and (x, y, z) , and t' and (x', y', z') , respectively. At time $t = t' = 0$ the origins O and O' and the x and x' axes of the two inertial systems coincide and at a later time t they have the relative location as depicted in Figure 5.1.

For convenience, let us introduce the two quantities

$$\beta = \frac{v}{c} \quad (5.1)$$

$$\gamma = \frac{1}{\sqrt{1 - \beta^2}} \quad (5.2)$$

where $v = |\mathbf{v}|$. In the following, we shall make frequent use of these shorthand notations.

As shown by Einstein, the two postulates of special relativity require that the spatial coordinates and times as measured by an observer in Σ and Σ' , respectively, are connected by the following transformation:

$$ct' = \gamma(ct - x\beta) \quad (5.3a)$$

$$x' = \gamma(x - vt) \quad (5.3b)$$

$$y' = y \quad (5.3c)$$

$$z' = z \quad (5.3d)$$

Taking the difference between the square of (5.3a) and the square of (5.3b) we find that

$$\begin{aligned} c^2t'^2 - x'^2 &= \gamma^2 (c^2t^2 - 2xc\beta t + x^2\beta^2 - x^2 + 2xvt - v^2t^2) \\ &= \frac{1}{1 - \frac{v^2}{c^2}} \left[c^2t^2 \left(1 - \frac{v^2}{c^2} \right) - x^2 \left(1 - \frac{v^2}{c^2} \right) \right] \\ &= c^2t^2 - x^2 \end{aligned} \quad (5.4)$$

From equation (5.3) on the facing page we see that the y and z coordinates are unaffected by the translational motion of the inertial system Σ' along the x axis of system Σ . Using this fact, we find that we can generalise the result in equation (5.4) above to

$$c^2t'^2 - x'^2 - y'^2 - z'^2 = c^2t^2 - x^2 - y^2 - z^2 \quad (5.5)$$

which means that if a light wave is transmitted from the coinciding origins O and O' at time $t = t' = 0$ it will arrive at an observer at (x, y, z) at time t in Σ and an observer at (x', y', z') at time t' in Σ' in such a way that both observers conclude that the speed (spatial distance divided by time) of light in vacuum is c . Hence, the speed of light in Σ and Σ' is the same. A linear coordinate transformation which has this property is called a (homogeneous) *Lorentz transformation*.

5.1.2 Lorentz space

Let us introduce an ordered quadruple of real numbers, enumerated with the help of upper indices $\mu = 0, 1, 2, 3$, where the zeroth component is ct (c is the speed of light and t is time), and the remaining components are the components of the

ordinary \mathbb{R}^3 radius vector \mathbf{x} defined in equation (M.1) on page 174:

$$x^\mu = (x^0, x^1, x^2, x^3) = (ct, x, y, z) \equiv (ct, \mathbf{x}) \quad (5.6)$$

We then interpret this quadruple x^μ as (the μ th component of) a *radius four-vector* in a real, linear, *four-dimensional vector space*.

Metric tensor

We want our space to be a *Riemannian space*, *i.e.*, a space where a distance and a scalar product are defined. We therefore need to define in this space a *metric tensor*, also known as the *fundamental tensor*, which we shall denote $g_{\mu\nu}$ and choose as (in matrix notation):

$$(g_{\mu\nu}) = \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & -1 & 0 & 0 \\ 0 & 0 & -1 & 0 \\ 0 & 0 & 0 & -1 \end{pmatrix} \quad (5.7)$$

i.e., with a main diagonal with sign sequence, or *signature*, $\{+, -, -, -\}$.¹

Radius four-vector in contravariant and covariant form

The radius four-vector $x^\mu = (x^0, x^1, x^2, x^3) = (ct, \mathbf{x})$, as defined in equation (5.6) above, is, by definition, the prototype of a *contravariant vector* (or, more accurately, a vector in *contravariant component form*). The corresponding *covariant vector* x_μ is obtained as (the upper index μ in x^μ is summed over and is therefore a *dummy index* and may be replaced by another dummy index ν):

$$x_\mu = g_{\mu\nu} x^\nu \quad (5.8)$$

This process is an example of *contraction* and is often called the “lowering” of index.

Index lowering of the contravariant radius four-vector x^μ amounts to multiplying the column vector representation of x^μ from the left by the matrix representation, equation (5.7), of $g_{\mu\nu}$ to obtain the column vector representation of x_μ . The simple diagonal form of $g_{\mu\nu}$, equation (5.7) above, means that the index lowering opera-

¹Without changing the physics, one can alternatively choose a signature $\{-, +, +, +\}$. The latter has the advantage that the transition from 3D to 4D becomes smooth, while it will introduce some annoying minus signs in the theory. In current physics literature, the signature $\{+, -, -, -\}$ seems to be the most commonly used one.

tion in our chosen flat 4D space is nearly trivial:

$$\begin{pmatrix} x_0 \\ x_1 \\ x_2 \\ x_3 \end{pmatrix} = \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & -1 & 0 & 0 \\ 0 & 0 & -1 & 0 \\ 0 & 0 & 0 & -1 \end{pmatrix} \begin{pmatrix} x^0 \\ x^1 \\ x^2 \\ x^3 \end{pmatrix} = \begin{pmatrix} x^0 \\ -x^1 \\ -x^2 \\ -x^3 \end{pmatrix} \quad (5.9)$$

which we can describe as

$$x_\mu = g_{\mu\nu}x^\nu = (ct, -\mathbf{x}) \quad (5.10)$$

i.e., the covariant radius four-vector x_μ is obtained from the contravariant radius four-vector x^μ simply by changing the sign of the last three components. These components are referred to as the *space components*; the zeroth component is referred to as the *time component*.

Scalar product and norm

Taking the scalar product of x^μ with itself, we get, by definition,

$$\begin{aligned} g_{\mu\nu}x^\nu x^\mu &= x_\mu x^\mu = (x_0, x_1, x_2, x_3) \cdot (x^0, x^1, x^2, x^3) = (ct, -\mathbf{x}) \cdot (ct, \mathbf{x}) \\ &= (ct, x, y, z) \cdot (ct, -x, -y, -z) = c^2t^2 - x^2 - y^2 - z^2 \end{aligned} \quad (5.11)$$

which acts as a *norm* or distance in our 4D space. We see, by comparing equation (5.11) and equation (5.5) on page 57, that this norm is conserved (invariant) during a Lorentz transformation. We notice further from equation (5.11) that our space has an *indefinite norm* which means that we deal with a *non-Euclidean space*. The four-dimensional space (or *space-time*) with these properties is called *Lorentz space* and is denoted \mathbb{L}^4 . The corresponding real, linear 4D space with a *positive definite norm* which is conserved during ordinary rotations is a *Euclidean vector space* which we denote \mathbb{R}^4 .

The \mathbb{L}^4 metric tensor equation (5.7) on the preceding page has a number of interesting properties: Firstly, we see that this tensor has a trace $\text{Tr}(g_{\mu\nu}) = -2$ whereas in \mathbb{R}^4 , as in any vector space with definite norm, the trace equals the space dimensionality. Secondly, we find, after trivial algebra, that the following relations between the contravariant, covariant and mixed forms of the metric tensor hold:

$$g_{\mu\nu} = g_{\nu\mu} \quad (5.12a)$$

$$g^{\mu\nu} = g_{\mu\nu} \quad (5.12b)$$

$$g_{\nu\kappa}g^{\kappa\mu} = g_\nu^\mu = \delta_\nu^\mu \quad (5.12c)$$

$$g^{\nu\kappa}g_{\kappa\mu} = g_\mu^\nu = \delta_\mu^\nu \quad (5.12d)$$

Here we have introduced the 4D version of the Kronecker delta δ_ν^μ , a mixed four-

tensor of rank 2 which fulfils

$$\delta_{\nu}^{\mu} = \delta_{\mu}^{\nu} = \begin{cases} 0 & \text{if } \mu \neq \nu \\ 1 & \text{if } \mu = \nu \end{cases} \quad (5.13)$$

Invariant line element and proper time

The *differential distance* ds between the two points x^{μ} and $x^{\mu} + dx^{\mu}$ in \mathbb{L}^4 can be calculated from the *Riemannian metric*, given by the *quadratic differential form*

$$\begin{aligned} ds^2 &= g_{\mu\nu} dx^{\nu} dx^{\mu} = dx_{\mu} dx^{\mu} = (dx^0)^2 - (dx^1)^2 - (dx^2)^2 - (dx^3)^2 \\ &= c^2 dt^2 - dx^2 - dy^2 - dz^2 \end{aligned} \quad (5.14)$$

where the metric tensor is as in equation (5.7) on page 58. The square root of this expression is the *invariant line element*

$$\begin{aligned} ds &= c dt \sqrt{1 - \frac{1}{c^2} \left[\left(\frac{dx}{dt} \right)^2 + \left(\frac{dy}{dt} \right)^2 + \left(\frac{dz}{dt} \right)^2 \right]} \\ &= c \sqrt{1 - \frac{v^2}{c^2}} dt = c \sqrt{1 - \beta^2} dt = \frac{c}{\gamma} dt = c d\tau \end{aligned} \quad (5.15)$$

where we introduced

$$d\tau = dt / \gamma \quad (5.16)$$

Since $d\tau$ measures the time when no spatial changes are present, it is called the *proper time*.

Expressing equation (5.5) on page 57 in terms of the differential interval ds and comparing with equation (5.14), we find that

$$ds^2 = c^2 dt^2 - dx^2 - dy^2 - dz^2 \quad (5.17)$$

is invariant during a Lorentz transformation. Conversely, we may say that every coordinate transformation which preserve this differential interval is a Lorentz transformation.

If in some inertial system

$$dx^2 + dy^2 + dz^2 < c^2 dt^2 \quad (5.18)$$

ds is a *time-like interval*, but if

$$dx^2 + dy^2 + dz^2 > c^2 dt^2 \quad (5.19)$$

ds is a *space-like interval*, whereas

$$dx^2 + dy^2 + dz^2 = c^2 dt^2 \quad (5.20)$$

is a *light-like interval*; we may also say that in this case we are on the *light cone*. A vector which has a light-like interval is called a *null vector*. The time-like,

space-like or light-like aspects of an interval ds is invariant under a Lorentz transformation.

Four-vector fields

Any quantity that relative to any coordinate system has a quadruple of real numbers and which transform in the same way as the radius four-vector x^μ , is called a *four-vector*. In analogy with the notation for the radius four-vector we introduce the notation $a^\mu = (a^0, \mathbf{a})$ for a general *contravariant four-vector field* in \mathbb{L}^4 and find that the “lowering of index” rule, equation (M.20) on page 178, for such an arbitrary four-vector yields the dual *covariant four-vector field*

$$a_\mu(x^\kappa) = g_{\mu\nu}a^\nu(x^\kappa) = (a^0(x^\kappa), -\mathbf{a}(x^\kappa)) \quad (5.21)$$

The scalar product between this four-vector field and another one $b^\mu(x^\kappa)$ is

$$g_{\mu\nu}a^\nu(x^\kappa)b^\mu(x^\kappa) = (a^0, -\mathbf{a}) \cdot (b^0, \mathbf{b}) = a^0b^0 - \mathbf{a} \cdot \mathbf{b} \quad (5.22)$$

which is a *scalar field*, *i.e.*, an invariant scalar quantity $\alpha(x^\kappa)$ which depends on time and space, as described by $x^\kappa = (ct, x, y, z)$.

The Lorentz transformation matrix

Introducing the transformation matrix

$$(\Lambda^\mu_\nu) = \begin{pmatrix} \gamma & -\beta\gamma & 0 & 0 \\ -\beta\gamma & \gamma & 0 & 0 \\ 0 & 0 & 1 & 0 \\ 0 & 0 & 0 & 1 \end{pmatrix} \quad (5.23)$$

the linear Lorentz transformation (5.3) on page 57, *i.e.*, the coordinate transformation $x^\mu \rightarrow x'^\mu = \Lambda'^\mu_\nu(x^0, x^1, x^2, x^3)$, from one inertial system Σ to another inertial system Σ' , can be written

$$x'^\mu = \Lambda^\mu_\nu x^\nu \quad (5.24)$$

The inverse transform then takes the form

$$x^\nu = (\Lambda^\mu_\nu)^{-1} x'^\mu \quad (5.25)$$

The Lorentz group

It is easy to show, by means of direct algebra, that two successive Lorentz transformations of the type in equation (5.25) above, and defined by the speed parameters

β_1 and β_2 , respectively, correspond to a single transformation with speed parameter

$$\beta = \frac{\beta_1 + \beta_2}{1 + \beta_1\beta_2} \quad (5.26)$$

This means that the nonempty set of Lorentz transformations constitute a *closed algebraic structure* with a binary operation which is *associative*. Furthermore, one can show that this set possesses at least one *identity element* and at least one *inverse element*. In other words, this set of Lorentz transformations constitute a *mathematical group*. However tempting, we shall not make any further use of *group theory*.

5.1.3 Minkowski space

Specifying a point $x^\mu = (x^0, x^1, x^2, x^3)$ in 4D space-time is a way of saying that “something takes place at a certain time $t = x^0/c$ and at a certain place $(x, y, z) = (x^1, x^2, x^3)$.” Such a point is therefore called an *event*. The trajectory for an event as a function of time and space is called a *world line*. For instance, the world line for a light ray which propagates in vacuum is the trajectory $x^0 = x^1$.

If we introduce

$$X^0 = ix^0 = ict \quad (5.27a)$$

$$dS = ids \quad (5.27b)$$

where $i = \sqrt{-1}$, we can rewrite equation (5.14) on page 60 as

$$dS^2 = (dX^0)^2 + (dx^1)^2 + (dx^2)^2 + (dx^3)^2 \quad (5.28)$$

i.e., as a 4D differential form which is *positive definite* just as is ordinary 3D *Euclidean space* \mathbb{R}^3 . We shall call the 4D Euclidean space constructed in this way the *Minkowski space* \mathbb{M}^4 .

As before, it suffices to consider the simplified case where the relative motion between Σ and Σ' is along the x axes. Then

$$dS^2 = (dX^0)^2 + (dx^1)^2 \quad (5.29)$$

and we consider X^0 and x^1 as orthogonal axes in an Euclidean space. As in all Euclidean spaces, every interval is invariant under a rotation of the X^0x^1 plane through an angle θ into $X'^0x'^1$:

$$X'^0 = -x^1 \sin \theta + X^0 \cos \theta \quad (5.30a)$$

$$x'^1 = x^1 \cos \theta + X^0 \sin \theta \quad (5.30b)$$

See Figure 5.2.

If we introduce the angle $\varphi = -i\theta$, often called the *rapidity* or the *Lorentz boost parameter*, and transform back to the original space and time variables by

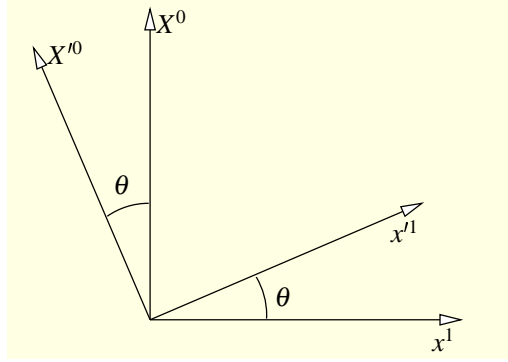


Figure 5.2. Minkowski space can be considered an ordinary Euclidean space where a Lorentz transformation from $(x^1, X^0 = ict)$ to $(x'^1, X'^0 = ict')$ corresponds to an ordinary rotation through an angle θ . This rotation leaves the Euclidean distance $(x^1)^2 + (X^0)^2 = x^2 - c^2t^2$ invariant.

using equation (5.27b) on the preceding page backwards, we obtain

$$ct' = -x \sinh \varphi + ct \cosh \varphi \quad (5.31a)$$

$$x' = x \cosh \varphi - ct \sinh \varphi \quad (5.31b)$$

which are identical to the transformation equations (5.3) on page 56 if we let

$$\sinh \varphi = \gamma\beta \quad (5.32a)$$

$$\cosh \varphi = \gamma \quad (5.32b)$$

$$\tanh \varphi = \beta \quad (5.32c)$$

It is therefore possible to envisage the Lorentz transformation as an “ordinary” rotation in the 4D Euclidean space \mathbb{M}^4 . equation (5.26) on the preceding page for successive Lorentz transformation then corresponds to the tanh addition formula

$$\tanh(\varphi_1 + \varphi_2) = \frac{\tanh \varphi_1 + \tanh \varphi_2}{1 + \tanh \varphi_1 \tanh \varphi_2} \quad (5.33)$$

The use of ict and \mathbb{M}^4 , which leads to the interpretation of the Lorentz transformation as an “ordinary” rotation, may, at best, be illustrative, but is not very physical. Besides, if we leave the flat \mathbb{L}^4 space and enter the curved space of general relativity, the “ ict ” trick will turn out to be an impasse. Let us therefore immediately return to \mathbb{L}^4 where all components are real valued.

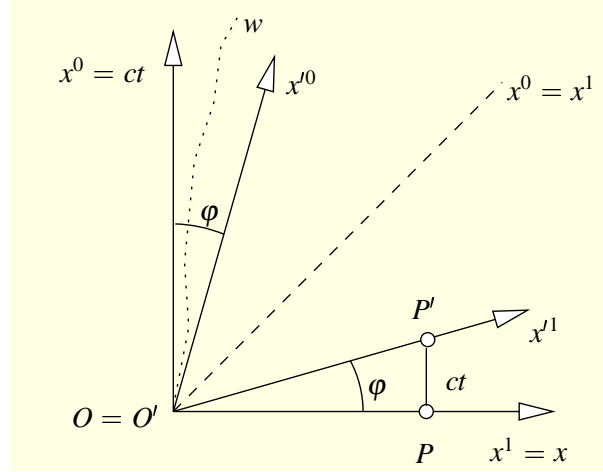


Figure 5.3. Minkowski diagram depicting geometrically the transformation (5.31) from the unprimed system to the primed system. Here w denotes the world line for an event and the line $x^0 = x^1 \Leftrightarrow x = ct$ the world line for a light ray in vacuum. Note that the event P is simultaneous with all points on the x^1 axis ($t = 0$), including the origin O while the event P' , which is also simultaneous with all points on the x' axis, including $O' = O$, to an observer at rest in the primed system, is not simultaneous with O in the unprimed system but occurs there at time $|P - P'|/c$.

5.2 Covariant classical mechanics

The measure ds of the differential “distance” in \mathbb{L}^4 allows us to define the *four-velocity*

$$u^\mu = \frac{dx^\mu}{ds} = \gamma \left(1, \frac{\mathbf{v}}{c} \right) = \left(\frac{1}{\sqrt{1 - \frac{v^2}{c^2}}}, \frac{\mathbf{v}}{c\sqrt{1 - \frac{v^2}{c^2}}} \right) \quad (5.34)$$

which, when multiplied with the scalar invariant m_0c^2 yields the *four-momentum*

$$p^\mu = m_0c^2 \frac{dx^\mu}{ds} = \gamma m_0c (c, \mathbf{v}) = \left(\frac{m_0c^2}{\sqrt{1 - \frac{v^2}{c^2}}}, \frac{m_0c\mathbf{v}}{\sqrt{1 - \frac{v^2}{c^2}}} \right) \quad (5.35)$$

From this we see that we can write

$$\mathbf{p} = m\mathbf{v} \quad (5.36)$$

where

$$m = \gamma m_0 = \frac{m_0}{\sqrt{1 - \frac{v^2}{c^2}}} \quad (5.37)$$

i.e., that Lorentz covariance implies that the mass-like term in the ordinary 3D linear momentum is not invariant.

The zeroth (time) component of the four-momentum p^μ is given by

$$p^0 = \gamma m_0 c^2 = \frac{m_0 c^2}{\sqrt{1 - \frac{v^2}{c^2}}} = mc^2 \quad (5.38)$$

We interpret this as the energy E , *i.e.*, we can write

$$p^\mu = (E, c\mathbf{p}) \quad (5.39)$$

Scalar multiplying this four-momentum with itself, we obtain

$$\begin{aligned} p_\mu p^\mu &= g_{\mu\nu} p^\nu p^\mu = (p^0)^2 - (p^1)^2 - (p^2)^2 - (p^3)^2 \\ &= (E, -c\mathbf{p}) \cdot (E, c\mathbf{p}) = E^2 - c^2 |\mathbf{p}|^2 \end{aligned} \quad (5.40)$$

$$= (m_0 c^2)^2 \frac{1 - \frac{v^2}{c^2}}{1 - \frac{v^2}{c^2}} = (m_0 c^2)^2 \quad (5.41)$$

Since this is an invariant, this equation holds in any inertial frame, particularly in the frame where $\mathbf{p} = 0$ where we thus have

$$E = m_0 c^2 \quad (5.42)$$

which is probably the most well-known physics formula ever.

5.3 Covariant classical electrodynamics

In the rest inertial system the charge density is ρ_0 . The four-vector (in contravariant component form)

$$j^\mu = \rho_0 \frac{dx^\mu}{ds} = \left(\rho, \rho \frac{\mathbf{v}}{c} \right) \quad (5.43)$$

where we introduced

$$\rho = \gamma \rho_0 \quad (5.44)$$

is called the *four-current*.

As is shown in example M.5 on page 185, the *d'Alembert operator* is the scalar

product of the four-del with itself:

$$\square^2 = \partial^\mu \partial_\mu = \partial_\mu \partial^\mu = \frac{1}{c^2} \frac{\partial^2}{\partial t^2} - \nabla^2 \quad (5.45)$$

Since it has the characteristics of a four-scalar, the d'Alembert operator is invariant and, hence, the homogeneous wave equation is Lorentz covariant.

5.3.1 The four-potential

If we introduce the *four-potential*

$$A^\mu = (\phi, c\mathbf{A}) \quad (5.46)$$

where ϕ is the scalar potential and \mathbf{A} the vector potential, defined in section 3.3 on page 34, we can write the inhomogeneous wave equations (Lorentz equations) equation (3.14) on page 36 in the following compact (and covariant) way:

$$\square^2 A^\mu = \frac{j^\mu}{\epsilon_0} \quad (5.47)$$

With the help of the above, we can formulate our electrodynamic equations covariantly. For instance, the covariant form of the *equation of continuity*, equation (1.21) on page 9 is

$$\partial_\mu j^\mu \equiv \frac{\partial j^\mu}{\partial x^\mu} = 0 \quad (5.48)$$

and the *Lorentz gauge condition*, equation (3.13) on page 36, can be written

$$\partial_\mu A^\mu \equiv \frac{\partial A^\mu}{\partial x^\mu} = 0 \quad (5.49)$$

The gauge transformations (3.15) on page 37 in covariant form are simply

$$A'^\mu = A^\mu - \partial^\mu [c\Gamma(x^\nu)] = A^\mu - \frac{\partial}{\partial x_\mu} c\Gamma(x^\nu) \quad (5.50)$$

If only one dimension Lorentz contracts (for instance, due to relative motion

along the x direction), a 3D spatial volume transforms according to

$$dV = d^3x = \frac{1}{\gamma} dV_0 = dV_0 \sqrt{1 - \beta^2} = dV_0 \sqrt{1 - \frac{v^2}{c^2}} \quad (5.51)$$

then from equation (5.44) on page 65 we see that

$$\rho dV = \rho_0 dV_0 \quad (5.52)$$

i.e., the charge in a given volume is conserved. We can therefore conclude that the electron charge is a universal constant.

5.3.2 The Liénard-Wiechert potentials

Let us now solve the Lorentz equation (the inhomogeneous wave equation) (3.14) on page 36 in vacuum for the case of a well-localised charge q' at a source point defined by the radius four-vector $x'^{\mu} = (x'^0 = ct', x'^1, x'^2, x'^3)$. The field (observation) point is denoted by the radius four-vector $x^{\mu} = (x^0 = ct, x^1, x^2, x^3)$.

In the rest system we know that the solution is simply

$$(A^{\mu})_0 = \left(\frac{q'}{4\pi\epsilon_0} \frac{1}{|\mathbf{x} - \mathbf{x}'|_0}, \mathbf{0} \right) \quad (5.53)$$

where $|\mathbf{x} - \mathbf{x}'|_0$ is the usual distance from the source point to the field point, evaluated in the rest system (signified by the index “0”).

Let us introduce the relative radius four-vector between the source point and the field point:

$$R^{\mu} = x^{\mu} - x'^{\mu} = (c(t - t'), \mathbf{x} - \mathbf{x}') \quad (5.54)$$

Scalar multiplying this relative four-vector with itself, we obtain

$$R^{\mu} R_{\mu} = (c(t - t'), \mathbf{x} - \mathbf{x}') \cdot (c(t - t'), -(\mathbf{x} - \mathbf{x}')) = c^2(t - t')^2 - |\mathbf{x} - \mathbf{x}'|^2 \quad (5.55)$$

We know that in vacuum the signal (field) from the charge q' at x'^{μ} propagates to x^{μ} with the speed of light c so that

$$|\mathbf{x} - \mathbf{x}'| = c(t - t') \quad (5.56)$$

Inserting this into equation (5.55), we see that

$$R^{\mu} R_{\mu} = 0 \quad (5.57)$$

or that equation (5.54) above can be written

$$R^{\mu} = (|\mathbf{x} - \mathbf{x}'|, \mathbf{x} - \mathbf{x}') \quad (5.58)$$

Now we want to find the correspondence to the rest system solution, equation (5.53), in an arbitrary inertial system. We note from equation (5.34) on page 64

that in the rest system

$$(u^\mu)_0 = \left(\frac{1}{\sqrt{1 - \frac{v^2}{c^2}}}, \frac{\mathbf{v}}{c\sqrt{1 - \frac{v^2}{c^2}}} \right)_{(\mathbf{v}=\mathbf{0})} = (1, \mathbf{0}) \quad (5.59)$$

and

$$(R^\mu)_0 = (|\mathbf{x} - \mathbf{x}'|, \mathbf{x} - \mathbf{x}')_0 = (|\mathbf{x} - \mathbf{x}'|_0, (\mathbf{x} - \mathbf{x}')_0) \quad (5.60)$$

Like all scalar products, $u^\mu R_\mu$ is invariant, so we can evaluate it in any inertial system. If we evaluate it in the rest system the result is:

$$\begin{aligned} u^\mu R_\mu &= (u^\mu R_\mu)_0 = (u^\mu)_0 (R_\mu)_0 \\ &= (1, \mathbf{0}) \cdot (|\mathbf{x} - \mathbf{x}'|_0, -(\mathbf{x} - \mathbf{x}')_0) = |\mathbf{x} - \mathbf{x}'|_0 \end{aligned} \quad (5.61)$$

We therefore see that the expression

$$A^\mu = \frac{q'}{4\pi\epsilon_0} \frac{u^\mu}{u^\nu R_\nu} \quad (5.62)$$

subject to the condition $R^\mu R_\mu = 0$ has the proper transformation properties (proper tensor form) and reduces, in the rest system, to the solution equation (5.53) on the preceding page. It is therefore the correct solution, valid in any inertial system.

According to equation (5.34) on page 64 and equation (5.58) on the preceding page

$$\begin{aligned} u^\nu R_\nu &= \gamma \left(1, \frac{\mathbf{v}}{c} \right) \cdot (|\mathbf{x} - \mathbf{x}'|, -(\mathbf{x} - \mathbf{x}')) \\ &= \gamma \left(|\mathbf{x} - \mathbf{x}'| - \frac{(\mathbf{x} - \mathbf{x}') \cdot \mathbf{v}}{c} \right) \end{aligned} \quad (5.63)$$

Introducing

$$s = |\mathbf{x} - \mathbf{x}'| - \frac{(\mathbf{x} - \mathbf{x}') \cdot \mathbf{v}}{c} \quad (5.64)$$

we can write

$$u^\nu R_\nu = \gamma s \quad (5.65)$$

and

$$\frac{u^\mu}{u^\nu R_\nu} = \left(\frac{1}{s}, \frac{\mathbf{v}}{cs} \right) \quad (5.66)$$

from which we see that the solution (5.62) can be written

$$A^\mu(x^\kappa) = \frac{q'}{4\pi\epsilon_0} \left(\frac{1}{s}, \frac{\mathbf{v}}{cs} \right) = (\phi, c\mathbf{A}) \quad (5.67)$$

where in the last step the definition of the four-potential, equation (5.46) on page 66, was used. Writing the solution in the ordinary 3D-way, we conclude that for a very

localised charge volume, moving relative an observer with a velocity \mathbf{v} , the scalar and vector potentials are given by the expressions

$$\phi(t, \mathbf{x}) = \frac{q'}{4\pi\epsilon_0} \frac{1}{s} = \frac{q'}{4\pi\epsilon_0} \frac{1}{|\mathbf{x} - \mathbf{x}'| - \frac{(\mathbf{x} - \mathbf{x}') \cdot \mathbf{v}}{c}} \quad (5.68a)$$

$$\mathbf{A}(t, \mathbf{x}) = \frac{q'}{4\pi\epsilon_0 c^2} \frac{\mathbf{v}}{s} = \frac{q'}{4\pi\epsilon_0 c^2} \frac{\mathbf{v}}{|\mathbf{x} - \mathbf{x}'| - \frac{(\mathbf{x} - \mathbf{x}') \cdot \mathbf{v}}{c}} \quad (5.68b)$$

These potentials are called the *Liénard-Wiechert potentials*.

5.3.3 The electromagnetic field tensor

Consider a vectorial (cross) product \mathbf{c} between two ordinary vectors \mathbf{a} and \mathbf{b} :

$$\begin{aligned} \mathbf{c} = \mathbf{a} \times \mathbf{b} &= \epsilon_{ijk} a_i b_j \hat{\mathbf{x}}_k \\ &= (a_2 b_3 - a_3 b_2) \hat{\mathbf{x}}_1 \\ &\quad + (a_3 b_1 - a_1 b_3) \hat{\mathbf{x}}_2 \\ &\quad + (a_1 b_2 - a_2 b_1) \hat{\mathbf{x}}_3 \end{aligned} \quad (5.69)$$

We notice that the k th component of the vector \mathbf{c} can be represented as

$$c_k = a_i b_j - a_j b_i = c_{ij} = -c_{ji}, \quad i, j \neq k \quad (5.70)$$

In other words, the *pseudovector* $\mathbf{c} = \mathbf{a} \times \mathbf{b}$ can be considered as an *antisymmetric tensor* of rank two!

The same is true for the curl operator $\nabla \times$. For instance, the Maxwell equation

$$\nabla \times \mathbf{E} = -\frac{\partial \mathbf{B}}{\partial t} \quad (5.71)$$

can in this tensor notation be written

$$\frac{\partial E_j}{\partial x^i} - \frac{\partial E_i}{\partial x^j} = -\frac{\partial B_{ij}}{\partial t} \quad (5.72)$$

We know from chapter 3 that the fields can be derived from the electromagnetic

potentials in the following way:

$$\mathbf{B} = \nabla \times \mathbf{A} \quad (5.73a)$$

$$\mathbf{E} = -\nabla\phi - \frac{\partial \mathbf{A}}{\partial t} \quad (5.73b)$$

In component form, this can be written

$$B_{ij} = \frac{\partial A_j}{\partial x^i} - \frac{\partial A_i}{\partial x^j} = \partial_i A_j - \partial_j A_i \quad (5.74a)$$

$$E_i = -\frac{\partial \phi}{\partial x^i} - \frac{\partial A_i}{\partial t} = -\partial_i \phi - \partial_t A_i \quad (5.74b)$$

From this, we notice the clear difference between the *axial vector* (pseudovector) \mathbf{B} and the *polar vector* (“ordinary vector”) \mathbf{E} .

Our goal is to express the electric and magnetic fields in a tensor form where the components are functions of the covariant form of the four-potential, equation (5.46) on page 66:

$$A_\mu = (\phi, -c\mathbf{A}) \quad (5.75)$$

Inspection of (5.75) and equation (5.74) above makes it natural to define the covariant four-tensor

$$F_{\mu\nu} = \frac{\partial A_\nu}{\partial x^\mu} - \frac{\partial A_\mu}{\partial x^\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu \quad (5.76)$$

This anti-symmetric (*skew-symmetric*), covariant four-tensor of rank 2 is called the *electromagnetic field tensor*. In matrix representation, the *covariant field tensor* can be written

$$F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu = \begin{pmatrix} 0 & E_x & E_y & E_z \\ -E_x & 0 & -cB_z & cB_y \\ -E_y & cB_z & 0 & -cB_x \\ -E_z & -cB_y & cB_x & 0 \end{pmatrix} \quad (5.77)$$

The matrix representation for the *contravariant field tensor* is

$$F^{\mu\nu} = \partial^\mu A^\nu - \partial^\nu A^\mu = \begin{pmatrix} 0 & -E_x & -E_y & -E_z \\ E_x & 0 & -cB_z & cB_y \\ E_y & cB_z & 0 & -cB_x \\ E_z & -cB_y & cB_x & 0 \end{pmatrix} \quad (5.78)$$

It is perhaps interesting to note that the field tensor is a sort of four-dimensional curl of the four-potential vector A^μ .

One can show that the two Maxwell source equations

$$\nabla \cdot \mathbf{E} = \frac{\rho}{\epsilon_0} \quad (5.79)$$

$$\nabla \times \mathbf{B} = \mu_0 \left(\mathbf{j} + \epsilon_0 \frac{\partial \mathbf{E}}{\partial t} \right) = \mu_0 \left(\rho \mathbf{v} + \epsilon_0 \frac{\partial \mathbf{E}}{\partial t} \right) \quad (5.80)$$

correspond to

$$\frac{\partial F^{\nu\mu}}{\partial x^\nu} = \frac{j^\mu}{\epsilon_0} \quad (5.81)$$

and that the two Maxwell “field” equations

$$\nabla \times \mathbf{E} = -\frac{\partial \mathbf{B}}{\partial t} \quad (5.82)$$

$$\nabla \cdot \mathbf{B} = 0 \quad (5.83)$$

correspond to

$$\frac{\partial F_{\mu\nu}}{\partial x^\kappa} + \frac{\partial F_{\nu\kappa}}{\partial x^\mu} + \frac{\partial F_{\kappa\mu}}{\partial x^\nu} = 0 \quad (5.84)$$

Hence, equation (5.81) and equation (5.84) above constitute Maxwell’s equations in four-dimensional formalism.

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Interactions of Fields and Particles

In this Chapter we study the interaction between electric and magnetic fields and electrically charged particles. The analysis is based on Lagrangian and Hamiltonian methods, is fully covariant, and yields results which are relativistically correct.

6.1 Charged Particles in an Electromagnetic Field

We first establish a relativistically correct theory describing the motion of charged particles in prescribed electric and magnetic fields. From these equations we may then calculate the charged particle dynamics in the most general case.

6.1.1 Covariant equations of motion

We can obtain an equation of motion if we find a Lagrange function in 4D for our problem and then apply a variational principle or if we find a Hamiltonian in 4D and solve the corresponding Hamilton's equations. We shall do both.

Lagrange formalism

Call the 4D Lagrange function $L_{(4)}$ and assume that it fulfils the variational principle

$$\delta \int_{x_0^\mu}^{x_1^\mu} L_{(4)}(x^\mu, u^\mu) ds = 0 \quad (6.1)$$

where ds is the invariant line element given by equation (5.15) on page 60, and the endpoints are fixed.

We must require that $L_{(4)}$ fulfils the following conditions:

1. The Lagrange function must be invariant. This implies that $L_{(4)}$ must be a scalar.
2. The Lagrange function must yield linear equations of motion. This implies that $L_{(4)}$ must not contain higher than the second power of the four-velocity u^μ .

According to formula (M.81) on page 189 the ordinary 3D Lagrangian is the difference between the kinetic and potential energies. A free particle has only kinetic energy. If the particle mass is m_0 then in 3D the kinetic energy is $m_0 v^2/2$. This suggests that in 4D the Lagrangian for a free particle should be

$$L_{(4)}^{\text{free}} = \frac{m_0 c^2}{2} u^\mu u_\mu \quad (6.2)$$

For an interaction with the electromagnetic field we can introduce the interaction with the help of the four-potential given by equation (5.75) on page 70 in the following way

$$L_{(4)} = \frac{m_0 c^2}{2} u^\mu u_\mu + q u_\mu A^\mu(x^\nu) \quad (6.3)$$

We call this the *four-Lagrangian*.

The variation principle (6.1) with the 4D Lagrangian (6.3) inserted, leads to

$$\begin{aligned} \delta \int_{x_0^\mu}^{x_1^\mu} L_{(4)}(x^\mu, u^\mu) ds &= \delta \int_{x_0^\mu}^{x_1^\mu} \left(\frac{m_0 c^2}{2} u^\mu u_\mu + q u^\mu A_\mu \right) ds = \\ &= \int_{x_0^\mu}^{x_1^\mu} \left[\frac{m_0 c^2}{2} \frac{\partial}{\partial u^\mu} (u^\mu u_\mu) \delta u^\mu + q \left(A_\mu \delta u^\mu + u^\mu \frac{\partial A_\mu}{\partial x^\nu} \delta x^\nu \right) \right] ds = \\ &= \int_{x_0^\mu}^{x_1^\mu} \left[m_0 c^2 u_\mu \delta u^\mu + q \left(A_\mu \delta u^\mu + u^\mu \frac{\partial A_\mu}{\partial x^\nu} \delta x^\nu \right) \right] ds = 0 \end{aligned} \quad (6.4)$$

According to equation (5.34) on page 64, the four-velocity is

$$u^\mu = \frac{dx^\mu}{ds} \quad (6.5)$$

which means that we can write the variation of u^μ as a total derivative with respect to s :

$$\delta u^\mu = \delta \left(\frac{dx^\mu}{ds} \right) = \frac{d}{ds} (\delta x^\mu) \quad (6.6)$$

Inserting this into the first two terms in the last integral in equation (6.4) on the

facing page, we obtain

$$\begin{aligned} & \delta \int_{x_0^\mu}^{x_1^\mu} L_{(4)}(x^\mu, u^\mu) ds \\ &= \int_{x_0^\mu}^{x_1^\mu} \left(m_0 c^2 u_\mu \frac{d}{ds} (\delta x^\mu) + q A_\mu \frac{d}{ds} (\delta x^\mu) + q u^\mu \frac{\partial A_\mu}{\partial x^\nu} \delta x^\nu \right) ds \end{aligned} \quad (6.7)$$

Partial integration in the two first terms in the right hand member of (6.7) gives

$$\begin{aligned} & \delta \int_{x_0^\mu}^{x_1^\mu} L_{(4)}(x^\mu, u^\mu) ds \\ &= \int_{x_0^\mu}^{x_1^\mu} \left(-m_0 c^2 \frac{du_\mu}{ds} \delta x^\mu - q \frac{dA_\mu}{ds} \delta x^\mu + q u^\mu \frac{\partial A_\mu}{\partial x^\nu} \delta x^\nu \right) ds \end{aligned} \quad (6.8)$$

where the integrated parts do not contribute since the variations at the endpoints vanishes. A change of irrelevant summation index from μ to ν in the first two terms of the right hand member of (6.8) yields, after moving the common factor δx^ν outside the partentthesis, the following expression:

$$\begin{aligned} & \delta \int_{x_0^\mu}^{x_1^\mu} L_{(4)}(x^\mu, u^\mu) ds \\ &= \int_{x_0^\mu}^{x_1^\mu} \left(-m_0 c^2 \frac{du_\nu}{ds} - q \frac{dA_\nu}{ds} + q u^\mu \frac{\partial A_\mu}{\partial x^\nu} \right) \delta x^\nu ds \end{aligned} \quad (6.9)$$

Applying well-known rules of differentiation and the expression (5.34) for the four-velocity, we can reform the expression for dA_ν/ds as follows:

$$\frac{dA_\nu}{ds} = \frac{\partial A_\nu}{\partial x^\mu} \frac{dx^\mu}{ds} = \frac{\partial A_\nu}{\partial x^\mu} u^\mu \quad (6.10)$$

By inserting this expression (6.10) into the right-hand member of equation (6.9) above, and moving the common factor $q u^\mu$ outside the bracket, we obtain the final variational principle expression

$$\begin{aligned} & \delta \int_{x_0^\mu}^{x_1^\mu} L_{(4)}(x^\mu, u^\mu) ds \\ &= \int_{x_0^\mu}^{x_1^\mu} \left[-m_0 c^2 \frac{du_\nu}{ds} + q u^\mu \left(\frac{\partial A_\mu}{\partial x^\nu} - \frac{\partial A_\nu}{\partial x^\mu} \right) \right] \delta x^\nu ds \end{aligned} \quad (6.11)$$

Since, according to the variational principle, this expression shall vanish and δx^ν is arbitrary along the world-line between the fixed end points x_0^μ and x_1^μ , the expression inside $[]$ in the integrand in the right hand member of (6.11) must vanish. In other words, we have found an equation of motion for a charged particle in a

prescribed electromagnetic field:

$$m_0 c^2 \frac{du_\nu}{ds} = qu^\mu \left(\frac{\partial A_\mu}{\partial x^\nu} - \frac{\partial A_\nu}{\partial x^\mu} \right) = qu^\mu (\partial_\nu A_\mu - \partial_\mu A_\nu) \quad (6.12)$$

With the help of equation (5.76) on page 70 we can express this equation in terms of the electromagnetic field tensor in the following way:

$$m_0 c^2 \frac{du_\nu}{ds} = qu^\mu F_{\nu\mu} \quad (6.13)$$

This is sought-for version of the covariant equation of motion for a particle in an electromagnetic field. It is often referred to as the *Minkowski equation*. As the reader can easily verify, the spatial part of this 4-vector equation is the covariant (relativistically correct) expression for the *Newton-Lorentz force equation*.

Hamiltonian formalism

The usual *Hamilton equations* for a 3D space are given by equation (M.86) on page 190 in chapter M. These six first-order partial differential equations are

$$\frac{\partial H}{\partial p_i} = \frac{dq_i}{dt} \quad (6.14a)$$

$$\frac{\partial H}{\partial q_i} = -\frac{dp_i}{dt} \quad (6.14b)$$

where $H(p_i, q_i, t) = p_i \dot{q}_i - L(q_i, \dot{q}_i, t)$ is the ordinary 3D Hamiltonian, q_i is a *generalised coordinate* and p_i is its *canonically conjugate momentum*.

We seek a similar set of equations in 4D space. To this end we introduce a *canonically conjugate four-momentum* p^μ in an analogous way as the ordinary 3D conjugate momentum:

$$p^\mu = \frac{\partial L_{(4)}}{\partial u_\mu} \quad (6.15)$$

and utilise the four-velocity u^μ , as given by equation (5.34) on page 64, to define the *four-Hamiltonian*

$$H_{(4)} = p^\mu u_\mu - L_{(4)} \quad (6.16)$$

With the help of these, the radius four-vector x^μ , considered as the *generalised four-coordinate*, and the invariant line element ds , defined in equation (5.15) on

page 60, we introduce the following eight partial differential equations:

$$\frac{\partial H_{(4)}}{\partial p^\mu} = \frac{dx_\mu}{ds} \quad (6.17a)$$

$$\frac{\partial H_{(4)}}{\partial x^\mu} = -\frac{dp_\mu}{ds} \quad (6.17b)$$

which form the *four-dimensional Hamilton equations*.

Our strategy now is to use equation (6.15) on the facing page and equations (6.17) above to derive an explicit algebraic expression for the canonically conjugate momentum four-vector. According to equation (5.39) on page 65, a four-momentum has a zeroth (time) component which we can identify with the total energy. Hence we require that the component p^0 of the conjugate four-momentum vector defined according to equation (6.15) on the preceding page be identical to the ordinary 3D Hamiltonian H and hence that this component solves the Hamilton equations equations (6.14) on the facing page. This later consistency check is left as an exercise to the reader.

Using the definition of $H_{(4)}$, equation (6.16) on the preceding page, and the expression for $L_{(4)}$, equation (6.3) on page 76, we obtain

$$H_{(4)} = p^\mu u_\mu - L_{(4)} = p^\mu u_\mu - \frac{m_0 c^2}{2} u^\mu u_\mu - q u^\mu A_\mu(x^\nu) \quad (6.18)$$

Furthermore, from the definition (6.15) of the conjugate four-momentum p^μ , we see that

$$\begin{aligned} p^\mu &= \frac{\partial L_{(4)}}{\partial u_\mu} = \frac{\partial}{\partial u_\mu} \left(\frac{m_0 c^2}{2} u^\mu u_\mu + q u^\mu A_\mu(x^\nu) \right) = \\ &= m_0 c^2 u^\mu + q A^\mu \end{aligned} \quad (6.19)$$

Inserting this into (6.18), we obtain

$$\begin{aligned} H_{(4)} &= m_0 c^2 u^\mu u_\mu + q A^\mu u_\mu - \frac{m_0 c^2}{2} u^\mu u_\mu - q u^\mu A_\mu(x^\nu) \\ &= \frac{m_0 c^2}{2} u^\mu u_\mu \end{aligned} \quad (6.20)$$

Since the four-velocity scalar-multiplied by itself is $u^\mu u_\mu = 1$, we clearly see from equation (6.20) above that $H_{(4)}$ is indeed a scalar invariant, whose value is

simply

$$H_{(4)} = \frac{m_0 c^2}{2} \quad (6.21)$$

However, at the same time (6.19) provides the algebraic relationship

$$u^\mu = \frac{1}{m_0 c^2} (p^\mu - q A^\mu) \quad (6.22)$$

and if this is used in (6.20) to eliminate u^μ , one gets

$$\begin{aligned} H_{(4)} &= \frac{m_0 c^2}{2} \left(\frac{1}{m_0 c^2} (p^\mu - q A^\mu) \frac{1}{m_0 c^2} (p_\mu - q A_\mu) \right) = \\ &= \frac{1}{2 m_0 c^2} (p^\mu - q A^\mu) (p_\mu - q A_\mu) = \\ &= \frac{1}{2 m_0 c^2} (p^\mu p_\mu - 2 q A^\mu p_\mu + q^2 A^\mu A_\mu) \end{aligned} \quad (6.23)$$

That this four-Hamiltonian yields the correct covariant equation of motion can be seen by inserting it into the four-dimensional Hamilton's equations (6.17) and using the relation (6.22):

$$\begin{aligned} \frac{\partial H_{(4)}}{\partial x^\mu} &= -\frac{q}{m_0 c^2} (p^\nu - q A^\nu) \frac{\partial A_\nu}{\partial x^\mu} \\ &= -\frac{q}{m_0 c^2} m_0 c^2 u^\nu \frac{\partial A_\nu}{\partial x^\mu} \\ &= -q u^\nu \frac{\partial A_\nu}{\partial x^\mu} \\ &= -\frac{d p_\mu}{d s} = -m_0 c^2 \frac{d u_\mu}{d s} - q \frac{\partial A_\mu}{\partial x^\nu} u^\nu \end{aligned} \quad (6.24)$$

where in the last step equation (6.19) on the preceding page was used. Rearranging terms, and using equation (5.78) on page 70, we obtain

$$m_0 c^2 \frac{d u_\mu}{d s} = q u^\nu \left(\frac{\partial A_\nu}{\partial x^\mu} - \frac{\partial A_\mu}{\partial x^\nu} \right) = q u^\nu F_{\mu\nu} \quad (6.25)$$

which is identical to the covariant equation of motion equation (6.13) on page 78. We can then safely conclude that the Hamiltonian in question is correct.

Using the fact that expression (6.23) above for $H_{(4)}$ is equal to the scalar value

$m_0c^2/2$, as derived in equation (6.21) on the facing page, the fact that

$$\begin{aligned} p^\mu &= (p^0, c\mathbf{p}) \\ A^\mu &= (\phi, c\mathbf{A}) \\ p^\mu p_\mu &= (p^0)^2 - c^2(\mathbf{p})^2 \\ A^\mu p_\mu &= \phi p^0 - c^2(\mathbf{p} \cdot \mathbf{A}) \\ A^\mu A_\mu &= \phi^2 - c^2(\mathbf{A})^2 \end{aligned}$$

and equation (6.19) on page 79, we obtain the equation

$$\frac{m_0c^2}{2} = \frac{1}{2m_0c^2} [(p^0)^2 - c^2(\mathbf{p})^2 - 2q\phi p^0 + 2qc^2(\mathbf{p} \cdot \mathbf{A}) + q^2\phi^2 - q^2c^2(\mathbf{A})^2] \quad (6.27)$$

which is a second order algebraic equation in p^0 which can be written

$$(m_0c^2)^2 = (p^0)^2 - 2q\phi p^0 - c^2(\mathbf{p})^2 + 2qc^2(\mathbf{p} \cdot \mathbf{A}) - q^2c^2(\mathbf{A})^2 + q^2\phi^2 \quad (6.28)$$

or

$$\begin{aligned} (p^0)^2 - 2q\phi p^0 - c^2 [(\mathbf{p})^2 - 2q\mathbf{p} \cdot \mathbf{A} + q^2(\mathbf{A})^2] + q^2\phi^2 - m_0^2c^4 \\ = (p^0)^2 - 2q\phi p^0 - c^2(\mathbf{p} - q\mathbf{A})^2 + q^2\phi^2 - m_0^2c^4 = 0 \end{aligned} \quad (6.29)$$

with two solution solutions

$$p^0 = q\phi \pm c\sqrt{(\mathbf{p} - q\mathbf{A})^2 + m_0^2c^2} \quad (6.30)$$

Since the fourth component (time component) p^0 of the four-momentum p^μ is the total energy, the positive solution in (6.30) must be identified with the ordinary Hamilton function H . This means that

$$H \equiv p^0 = q\phi + c\sqrt{(\mathbf{p} - q\mathbf{A})^2 + m_0^2c^2} \quad (6.31)$$

is the ordinary 3D Hamilton function for a charged particle moving in scalar and vector potentials associated with prescribed electric and magnetic fields.

The ordinary Lagrange and Hamilton functions L and H are related to each other by the 3D transformation [*cf.* the 4D transformation (6.16) between $L_{(4)}$ and $H_{(4)}$]

$$L = \mathbf{p} \cdot \mathbf{v} - H \quad (6.32)$$

Using the explicit expressions (equation (6.31)) and (equation (6.32) above), we

obtain the explicit expression for the ordinary 3D Lagrange function

$$L = \mathbf{p} \cdot \mathbf{v} - q\phi - c\sqrt{(\mathbf{p} - q\mathbf{A})^2 + m_0^2 c^2} \quad (6.33)$$

and if we make the identification

$$\mathbf{p} - q\mathbf{A} = \frac{m_0 \mathbf{v}}{\sqrt{1 - \frac{v^2}{c^2}}} = m\mathbf{v} \quad (6.34)$$

where the quantity $m\mathbf{v}$ is called the *kinetic momentum*, we can rewrite this expression for the ordinary Lagrangian as follows:

$$\begin{aligned} L &= q\mathbf{A} \cdot \mathbf{v} + mv^2 - q\phi - c\sqrt{m^2 v^2 - m_0^2 c^2} \\ &= -q\phi + q\mathbf{A} \cdot \mathbf{v} - m_0 c^2 \sqrt{1 - \frac{v^2}{c^2}} \end{aligned} \quad (6.35)$$

What we have obtained is the correct expression for the Lagrangian describing the motion of a charged particle in scalar and vector potentials associated with prescribed electric and magnetic fields.

6.2 Covariant Field Theory

So far, we have considered two classes of problems. Either we have calculated the fields from given, prescribed distributions of charges and currents, or we have derived the equations of motion for charged particles in given, prescribed fields. Let us now put the fields and the particles on an equal footing and present a theoretical description which treats the fields, the particles, and their interactions in a unified way. This involves transition to a field picture with an infinite number of degrees of freedom. We shall first consider a simple mechanical problem whose solution is well known. Then, drawing inferences from this model problem, we apply a similar view on the electromagnetic problem.

6.2.1 Lagrange-Hamilton formalism for fields and interactions

Consider N identical mass points, each with mass m and connected to its neighbour along a one-dimensional straight line, which we choose to be the x axis, by identical ideal springs with spring constants k . At equilibrium the mass points are at rest, distributed evenly with a distance a to their two nearest neighbours. After

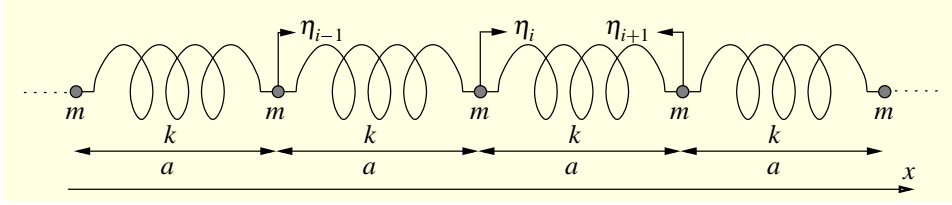


Figure 6.1. A one-dimensional chain consisting of N discrete, identical mass points m , connected to their neighbours with identical, ideal springs with spring constants k . The equilibrium distance between the neighbouring mass points is a and $\eta_{i-1}(t)$, $\eta_i(t)$, $\eta_{i+1}(t)$ are the instantaneous deviations, along the x axis, of positions of the $(i-1)$ th, i th, and $(i+1)$ th mass point, respectively.

perturbation, the motion of mass point i will be a one-dimensional oscillatory motion along \hat{x} . Let us denote the magnitude of the deviation for mass point i from its equilibrium position by $\eta_i(t)\hat{x}$.

The solution to this mechanical problem can be obtained if we can find a *Lagrangian* (*Lagrange function*) L which satisfies the variational equation

$$\delta \int L(\eta_i, \dot{\eta}_i, t) dt = 0 \quad (6.36)$$

According to equation (M.81) on page 189, the Lagrangian is $L = T - V$ where T denotes the *kinetic energy* and V the *potential energy* of a classical mechanical system with *conservative forces*. In our case the Lagrangian is

$$L = \frac{1}{2} \sum_{i=1}^N [m\dot{\eta}_i^2 - k(\eta_{i+1} - \eta_i)^2] \quad (6.37)$$

Let us write the Lagrangian, as given by equation (6.37), in the following way:

$$L = \sum_{i=1}^N a\mathcal{L}_i \quad (6.38)$$

Here,

$$\mathcal{L}_i = \frac{1}{2} \left[\frac{m}{a} \dot{\eta}_i^2 - ka \left(\frac{\eta_{i+1} - \eta_i}{a} \right)^2 \right] \quad (6.39)$$

is the so called linear *Lagrange density*. If we now let $N \rightarrow \infty$ and, at the same time,

let the springs become infinitesimally short according to the following scheme:

$$a \rightarrow dx \quad (6.40a)$$

$$\frac{m}{a} \rightarrow \frac{dm}{dx} = \mu \quad \text{linear mass density} \quad (6.40b)$$

$$ka \rightarrow Y \quad \text{Young's modulus} \quad (6.40c)$$

$$\frac{\eta_{i+1} - \eta_i}{a} \rightarrow \frac{\partial \eta}{\partial x} \quad (6.40d)$$

we obtain

$$L = \int \mathcal{L} dx \quad (6.41)$$

where

$$\mathcal{L} \left(\eta, \frac{\partial \eta}{\partial t}, \frac{\partial \eta}{\partial x}, t \right) = \frac{1}{2} \left[\mu \left(\frac{\partial \eta}{\partial t} \right)^2 - Y \left(\frac{\partial \eta}{\partial x} \right)^2 \right] \quad (6.42)$$

Notice how we made a transition from a discrete description, in which the mass points were identified by a discrete integer variable $i = 1, 2, \dots, N$, to a continuous description, where the infinitesimal mass points were instead identified by a continuous real parameter x , namely their position along \hat{x} .

A consequence of this transition is that the number of degrees of freedom for the system went from the finite number N to infinity! Another consequence is that \mathcal{L} has now become dependent also on the partial derivative with respect to x of the “field coordinate” η . But, as we shall see, the transition is well worth the price because it allows us to treat all fields, be it classical scalar or vectorial fields, or wave functions, spinors and other fields that appear in quantum physics, on an equal footing.

Under the assumption of time independence and fixed endpoints, the variation principle (6.36) on the previous page yields:

$$\begin{aligned} & \delta \int L dt \\ &= \delta \iint \mathcal{L} \left(\eta, \frac{\partial \eta}{\partial t}, \frac{\partial \eta}{\partial x} \right) dx dt \\ &= \iint \left[\frac{\partial \mathcal{L}}{\partial \eta} \delta \eta + \frac{\partial \mathcal{L}}{\partial \left(\frac{\partial \eta}{\partial t} \right)} \delta \left(\frac{\partial \eta}{\partial t} \right) + \frac{\partial \mathcal{L}}{\partial \left(\frac{\partial \eta}{\partial x} \right)} \delta \left(\frac{\partial \eta}{\partial x} \right) \right] dx dt \\ &= 0 \end{aligned} \quad (6.43)$$

The last integral can be integrated by parts. This results in the expression

$$\iint \left[\frac{\partial \mathcal{L}}{\partial \eta} - \frac{\partial}{\partial t} \left(\frac{\partial \mathcal{L}}{\partial \left(\frac{\partial \eta}{\partial t} \right)} \right) - \frac{\partial}{\partial x} \left(\frac{\partial \mathcal{L}}{\partial \left(\frac{\partial \eta}{\partial x} \right)} \right) \right] \delta \eta \, dx \, dt = 0 \quad (6.44)$$

where the variation is arbitrary (and the endpoints fixed). This means that the integrand itself must vanish. If we introduce the *functional derivative*

$$\frac{\delta \mathcal{L}}{\delta \eta} = \frac{\partial \mathcal{L}}{\partial \eta} - \frac{\partial}{\partial x} \left(\frac{\partial \mathcal{L}}{\partial \left(\frac{\partial \eta}{\partial x} \right)} \right) \quad (6.45)$$

we can express this as

$$\frac{\delta \mathcal{L}}{\delta \eta} - \frac{\partial}{\partial t} \left(\frac{\partial \mathcal{L}}{\partial \left(\frac{\partial \eta}{\partial t} \right)} \right) = 0 \quad (6.46)$$

which is the one-dimensional *Euler-Lagrange equation*.

Inserting the linear mass point chain Lagrangian density, equation (6.42) on the facing page, into equation (6.46) above, we obtain the equation of motion for our one-dimensional linear mechanical structure. It is:

$$\mu \frac{\partial^2}{\partial t^2} \eta - Y \frac{\partial^2}{\partial x^2} \eta = \left(\frac{\mu}{Y} \frac{\partial^2}{\partial t^2} - \frac{\partial^2}{\partial x^2} \right) \eta = 0 \quad (6.47)$$

i.e., the one-dimensional wave equation for compression waves which propagate with phase speed $v_\phi = \sqrt{Y/\mu}$ along the linear structure.

A generalisation of the above 1D results to a three-dimensional continuum is straightforward. For this 3D case we get the variational principle

$$\begin{aligned} \delta \int L \, dt &= \delta \int \mathcal{L} \, d^3x \, dt = \delta \iint \mathcal{L} \left(\eta, \frac{\partial \eta}{\partial x^\mu} \right) d^4x \\ &= \iint \left[\frac{\partial \mathcal{L}}{\partial \eta} - \frac{\partial}{\partial x^\mu} \left(\frac{\partial \mathcal{L}}{\partial \left(\frac{\partial \eta}{\partial x^\mu} \right)} \right) \right] \delta \eta \, d^4x = 0 \end{aligned} \quad (6.48)$$

where the variation $\delta \eta$ is arbitrary and the endpoints are fixed. This means that the

integrand itself must vanish:

$$\frac{\partial \mathcal{L}}{\partial \eta} - \frac{\partial}{\partial x^\mu} \left(\frac{\partial \mathcal{L}}{\partial \left(\frac{\partial \eta}{\partial x^\mu} \right)} \right) = 0 \quad (6.49)$$

This constitutes the three-dimensional *Euler-Lagrange equations*.

Introducing the *three-dimensional functional derivative*

$$\frac{\delta \mathcal{L}}{\delta \eta} = \frac{\partial \mathcal{L}}{\partial \eta} - \frac{\partial}{\partial x^i} \left(\frac{\partial \mathcal{L}}{\partial \left(\frac{\partial \eta}{\partial x^i} \right)} \right) \quad (6.50)$$

we can express this as

$$\frac{\delta \mathcal{L}}{\delta \eta} - \frac{\partial}{\partial t} \left(\frac{\partial \mathcal{L}}{\partial \left(\frac{\partial \eta}{\partial t} \right)} \right) = 0 \quad (6.51)$$

In analogy with particle mechanics (finite number of degrees of freedom), we may introduce the *canonically conjugate momentum density*

$$\pi(x^\mu) = \pi(t, \mathbf{x}) = \frac{\partial \mathcal{L}}{\partial \left(\frac{\partial \eta}{\partial t} \right)} \quad (6.52)$$

and define the *Hamilton density*

$$\mathcal{H} \left(\pi, \eta, \frac{\partial \eta}{\partial x^i}; t \right) = \pi \frac{\partial \eta}{\partial t} - \mathcal{L} \left(\eta, \frac{\partial \eta}{\partial t}, \frac{\partial \eta}{\partial x^i} \right) \quad (6.53)$$

If, as usual, we differentiate this expression and identify terms, we obtain the following *Hamilton density equations*

$$\frac{\partial \mathcal{H}}{\partial \pi} = \frac{\partial \eta}{\partial t} \quad (6.54a)$$

$$\frac{\delta \mathcal{H}}{\delta \eta} = - \frac{\partial \pi}{\partial t} \quad (6.54b)$$

The Hamilton density functions are in many ways similar to the ordinary Hamilton functions and lead to similar results.

The electromagnetic field

Above, when we described the mechanical field, we used a scalar field $\eta(t, \mathbf{x})$. If we want to describe the electromagnetic field in terms of a Lagrange density \mathcal{L} and Euler-Lagrange equations, it comes natural to express \mathcal{L} in terms of the four-potential $A^\mu(x^\kappa)$.

The entire system of particles and fields consists of a mechanical part, a field part and an interaction part. We therefore assume that the total Lagrange density

\mathcal{L}^{tot} for this system can be expressed as

$$\mathcal{L}^{\text{tot}} = \mathcal{L}^{\text{mech}} + \mathcal{L}^{\text{inter}} + \mathcal{L}^{\text{field}} \quad (6.55)$$

where the mechanical part has to do with the particle motion (kinetic energy). It is given by $L_{(4)}/V$ where $L_{(4)}$ is given by equation (6.2) on page 76 and V is the volume. Expressed in the *rest mass density* ρ_m^0 , the *mechanical Lagrange density* can be written

$$\mathcal{L}^{\text{mech}} = \frac{1}{2}\rho_m^0 c^2 u^\mu u_\mu \quad (6.56)$$

The $\mathcal{L}^{\text{inter}}$ part which describes the interaction between the charged particles and the external electromagnetic field. A convenient expression for this *interaction Lagrange density* is

$$\mathcal{L}^{\text{inter}} = j^\mu A_\mu \quad (6.57)$$

For the field part $\mathcal{L}^{\text{field}}$ we choose the difference between magnetic and electric energy density (in analogy with the difference between kinetic and potential energy in a mechanical field). Using the field tensor, we express this *field Lagrange density* as

$$\mathcal{L}^{\text{field}} = \frac{1}{4}\epsilon_0 F^{\mu\nu} F_{\mu\nu} \quad (6.58)$$

so that the total Lagrangian density can be written

$$\mathcal{L}^{\text{tot}} = \frac{1}{2}\rho_m^0 c^2 u^\mu u_\mu + j^\mu A_\mu + \frac{1}{4}\epsilon_0 F^{\mu\nu} F_{\mu\nu} \quad (6.59)$$

▷FIELD ENERGY DIFFERENCE EXPRESSED IN THE FIELD TENSOR

EXAMPLE 6.1

Show, by explicit calculation, that

$$F^{\mu\nu} F_{\mu\nu} = 2(c^2 B^2 - E^2) \quad (6.60)$$

From formula (5.78) on page 70 we recall that

$$F^{\mu\nu} = \partial^\mu A^\nu - \partial^\nu A^\mu = \begin{pmatrix} 0 & -E_x & -E_y & -E_z \\ E_x & 0 & -cB_z & cB_y \\ E_y & cB_z & 0 & -cB_x \\ E_z & -cB_y & cB_x & 0 \end{pmatrix} \quad (6.61)$$

and from formula (5.77) on page 70 that

$$F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu = \begin{pmatrix} 0 & E_x & E_y & E_z \\ -E_x & 0 & -cB_z & cB_y \\ -E_y & cB_z & 0 & -cB_x \\ -E_z & -cB_y & cB_x & 0 \end{pmatrix} \quad (6.62)$$

where μ denotes the row number and ν the column number. Then, Einstein summation

and direct substitution yields

$$\begin{aligned}
F^{\mu\nu}F_{\mu\nu} &= F^{00}F_{00} + F^{01}F_{01} + F^{02}F_{02} + F^{03}F_{03} \\
&\quad + F^{10}F_{10} + F^{11}F_{11} + F^{12}F_{12} + F^{13}F_{13} \\
&\quad + F^{20}F_{20} + F^{21}F_{21} + F^{22}F_{22} + F^{23}F_{23} \\
&\quad + F^{30}F_{30} + F^{31}F_{31} + F^{32}F_{32} + F^{33}F_{33} \\
&= 0 - E_x^2 - E_y^2 - E_z^2 \\
&\quad - E_x^2 + 0 + c^2B_z^2 + c^2B_y^2 \\
&\quad - E_y^2 + c^2B_z^2 + 0 + c^2B_x^2 \\
&\quad - E_z^2 + c^2B_y^2 + c^2B_x^2 + 0 \\
&= -2E_x^2 - 2E_y^2 - 2E_z^2 + c^2B_x^2 + c^2B_y^2 + c^2B_z^2 \\
&= -2E^2 + 2c^2B^2 = 2(c^2B^2 - E^2)
\end{aligned} \tag{6.63}$$

QED■

END OF EXAMPLE 6.1◀

Using \mathcal{L}^{tot} in the 3D Euler-Lagrange equations, equation (6.49) on page 86 (with η replaced by A_ν), we can derive the dynamics for the whole system. For instance, the electromagnetic part of the Lagrangian density

$$\mathcal{L}^{\text{EM}} = \mathcal{L}^{\text{inter}} + \mathcal{L}^{\text{field}} = j^\nu A_\nu + \frac{1}{4}\epsilon_0 F^{\mu\nu}F_{\mu\nu} \tag{6.64}$$

inserted into the Euler-Lagrange equations, expression (6.49) on page 86, yields two of Maxwell's equations. To see this, we note from equation (6.64) above and the results in Example 6.1 that

$$\frac{\partial \mathcal{L}^{\text{EM}}}{\partial A_\nu} = j^\nu \tag{6.65}$$

Furthermore,

$$\begin{aligned}
\partial_\mu \left[\frac{\partial \mathcal{L}^{\text{EM}}}{\partial(\partial_\mu A_\nu)} \right] &= \frac{\epsilon_0}{4} \partial_\mu \left[\frac{\partial}{\partial(\partial_\mu A_\nu)} \left(F^{\kappa\lambda} F_{\kappa\lambda} \right) \right] \\
&= \frac{\epsilon_0}{4} \partial_\mu \left\{ \frac{\partial}{\partial(\partial_\mu A_\nu)} \left[(\partial^\kappa A^\lambda - \partial^\lambda A^\kappa)(\partial_\kappa A_\lambda - \partial_\lambda A_\kappa) \right] \right\} \\
&= \frac{\epsilon_0}{4} \partial_\mu \left\{ \frac{\partial}{\partial(\partial_\mu A_\nu)} \left[\partial^\kappa A^\lambda \partial_\kappa A_\lambda - \partial^\kappa A^\lambda \partial_\lambda A_\kappa \right. \right. \\
&\quad \left. \left. - \partial^\lambda A^\kappa \partial_\kappa A_\lambda + \partial^\lambda A^\kappa \partial_\lambda A_\kappa \right] \right\} \\
&= \frac{\epsilon_0}{2} \partial_\mu \left[\frac{\partial}{\partial(\partial_\mu A_\nu)} \left(\partial^\kappa A^\lambda \partial_\kappa A_\lambda - \partial^\kappa A^\lambda \partial_\lambda A_\kappa \right) \right]
\end{aligned} \tag{6.66}$$

But

$$\begin{aligned}
\frac{\partial}{\partial(\partial_\mu A_\nu)} \left(\partial^\kappa A^\lambda \partial_\kappa A_\lambda \right) &= \partial^\kappa A^\lambda \frac{\partial}{\partial(\partial_\mu A_\nu)} \partial_\kappa A_\lambda \\
&\quad + \partial_\kappa A_\lambda \frac{\partial}{\partial(\partial_\mu A_\nu)} \partial^\kappa A^\lambda \\
&= \partial^\kappa A^\lambda \frac{\partial}{\partial(\partial_\mu A_\nu)} \partial_\kappa A_\lambda \\
&\quad + \partial_\kappa A_\lambda \frac{\partial}{\partial(\partial_\mu A_\nu)} g^{\kappa\alpha} \partial_\alpha g^{\lambda\beta} A_\beta \\
&= \partial^\kappa A^\lambda \frac{\partial}{\partial(\partial_\mu A_\nu)} \partial_\kappa A_\lambda \\
&\quad + g^{\kappa\alpha} g^{\lambda\beta} \partial_\kappa A_\lambda \frac{\partial}{\partial(\partial_\mu A_\nu)} \partial_\alpha A_\beta \\
&= \partial^\kappa A^\lambda \frac{\partial}{\partial(\partial_\mu A_\nu)} \partial_\kappa A_\lambda \\
&\quad + \partial^\alpha A^\beta \frac{\partial}{\partial(\partial_\mu A_\nu)} \partial_\alpha A_\beta \\
&= 2\partial^\mu A^\nu
\end{aligned} \tag{6.67}$$

Similarly,

$$\frac{\partial}{\partial(\partial_\mu A_\nu)} \left(\partial^\kappa A^\lambda \partial_\lambda A_\kappa \right) = 2\partial^\nu A^\mu \tag{6.68}$$

so that

$$\partial_\mu \left[\frac{\partial \mathcal{L}^{\text{EM}}}{\partial (\partial_\mu A_\nu)} \right] = \epsilon_0 \partial_\mu (\partial^\mu A^\nu - \partial^\nu A^\mu) = \epsilon_0 \frac{\partial F^{\mu\nu}}{\partial x^\mu} \quad (6.69)$$

This means that the Euler-Lagrange equations, expression (6.49) on page 86, for the Lagrangian density \mathcal{L}^{EM} and with A_ν as the field quantity become

$$\frac{\partial \mathcal{L}^{\text{EM}}}{\partial A_\nu} - \partial_\mu \left[\frac{\partial \mathcal{L}^{\text{EM}}}{\partial (\partial_\mu A_\nu)} \right] = j^\nu - \epsilon_0 \frac{\partial F^{\mu\nu}}{\partial x^\mu} = 0 \quad (6.70)$$

or

$$\frac{\partial F^{\mu\nu}}{\partial x^\mu} = \frac{j^\nu}{\epsilon_0} \quad (6.71)$$

Explicitly, setting $\nu = 0$ in this covariant equation and using the matrix representation formula (5.78) on page 70 for the covariant component form of the electromagnetic field tensor $F^{\mu\nu}$, we obtain

$$\begin{aligned} \frac{\partial F^{00}}{\partial x^0} + \frac{\partial F^{10}}{\partial x^1} + \frac{\partial F^{20}}{\partial x^2} + \frac{\partial F^{30}}{\partial x^3} &= 0 + \frac{\partial E_x}{\partial x} + \frac{\partial E_y}{\partial y} + \frac{\partial E_z}{\partial z} \\ &= \nabla \cdot \mathbf{E} = \frac{\rho}{\epsilon_0} \end{aligned} \quad (6.72)$$

which is the Maxwell source equation for the electric field, equation (1.43a) on page 14. For $\nu = 1$ we get

$$\begin{aligned} \frac{\partial F^{01}}{\partial x^0} + \frac{\partial F^{11}}{\partial x^1} + \frac{\partial F^{21}}{\partial x^2} + \frac{\partial F^{31}}{\partial x^3} &= -\frac{1}{c} \frac{\partial E_x}{\partial t} + 0 - c \frac{\partial B_z}{\partial y} + c \frac{\partial B_y}{\partial z} \\ &= \frac{\rho u_x}{\epsilon_0} \end{aligned} \quad (6.73)$$

or, using $\epsilon_0 \mu_0 = 1/c^2$ and identifying $\rho u_x = j_x$,

$$\frac{\partial B_y}{\partial z} - \frac{\partial B_z}{\partial y} - \epsilon_0 \mu_0 \frac{\partial E_x}{\partial t} = \mu_0 j_x \quad (6.74)$$

and similarly for $\nu = 2, 3$. In summary, in three-vector form, we can write the

result as

$$\nabla \times \mathbf{B} - \epsilon_0 \mu_0 \frac{\partial \mathbf{E}}{\partial t} = \mu_0 \mathbf{j}(t, \mathbf{x}) \quad (6.75)$$

which is the Maxwell source equation for the magnetic field, equation (1.43d) on page 14.

Other fields

In general, the dynamic equations for most any fields, and not only electromagnetic ones, can be derived from a Lagrangian density together with a variational principle (the Euler-Lagrange equations). Both linear and non-linear fields are studied with this technique. As a simple example, consider a real, scalar field η which has the following Lagrange density:

$$\mathcal{L} = \frac{1}{2} (\partial_\mu \eta \partial^\mu \eta - m^2 \eta^2) = \frac{1}{2} \left(\frac{\partial \eta}{\partial x^\mu} \frac{\partial \eta}{\partial x_\mu} - m^2 \eta^2 \right) \quad (6.76)$$

Insertion into the 1D Euler-Lagrange equation, equation (6.46) on page 85, yields the dynamic equation

$$(\square^2 - m^2)\eta = 0 \quad (6.77)$$

with the solution

$$\eta = e^{i(\mathbf{k} \cdot \mathbf{x} - \omega t)} \frac{e^{-m|\mathbf{x}|}}{|\mathbf{x}|} \quad (6.78)$$

which describes the *Yukawa meson field* for a scalar meson with mass m . With

$$\pi = \frac{1}{c^2} \frac{\partial \eta}{\partial t} \quad (6.79)$$

we obtain the Hamilton density

$$\mathcal{H} = \frac{1}{2} \left[c^2 \pi^2 + (\nabla \eta)^2 + m^2 \eta^2 \right] \quad (6.80)$$

which is positive definite.

Another Lagrangian density which has attracted quite some interest is the *Proca*

Lagrangian

$$\mathcal{L}^{\text{EM}} = \mathcal{L}^{\text{inter}} + \mathcal{L}^{\text{field}} = j^\nu A_\nu + \frac{1}{4} \epsilon_0 F^{\mu\nu} F_{\mu\nu} + m^2 A^\mu A_\mu \quad (6.81)$$

which leads to the dynamic equation

$$\frac{\partial F^{\mu\nu}}{\partial x^\mu} + m^2 A^\nu = \frac{j^\nu}{\epsilon_0} \quad (6.82)$$

This equation describes an electromagnetic field with a mass, or, in other words, *massive photons*. If massive photons would exist, large-scale magnetic fields, including those of the earth and galactic spiral arms, would be significantly modified to yield measurable discrepancies from their usual form. Space experiments of this kind onboard satellites have led to stringent upper bounds on the photon mass. If the photon really has a mass, it will have an impact on electrodynamics as well as on cosmology and astrophysics.

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Interactions of Fields and Matter

The microscopic Maxwell equations (1.43) derived in chapter 1 are valid on all scales where a classical description is good. However, when macroscopic matter is present, it is sometimes convenient to use the corresponding macroscopic Maxwell equations (in a statistical sense) in which auxiliary, derived fields are introduced in order to incorporate effects of macroscopic matter when this is immersed fully or partially in an electromagnetic field.

7.1 Electric polarisation and the electric displacement vector

7.1.1 Electric multipole moments

The electrostatic properties of a spatial volume containing electric charges and located near a point \mathbf{x}_0 can be characterized in terms of the *total charge* or *electric monopole moment*

$$q = \int_V \rho(\mathbf{x}') d^3x' \quad (7.1)$$

where the ρ is the charge density introduced in equation (1.7) on page 4), the *electric dipole moment vector*

$$\mathbf{p} = \int_V (\mathbf{x}' - \mathbf{x}_0) \rho(\mathbf{x}') d^3x' \quad (7.2)$$

with components p_i , $i = 1, 2, 3$, the *electric quadrupole moment tensor*

$$\mathbf{Q} = \int_V (\mathbf{x}' - \mathbf{x}_0)(\mathbf{x}' - \mathbf{x}_0) \rho(\mathbf{x}') d^3x' \quad (7.3)$$

with components Q_{ij} , $i, j = 1, 2, 3$, and higher order electric moments.

In particular, the electrostatic potential equation (3.3) on page 33 from a charge distribution located near \mathbf{x}_0 can be Taylor expanded in the following way:

$$\phi^{\text{stat}}(\mathbf{x}) = \frac{1}{4\pi\epsilon_0} \left[\frac{q}{|\mathbf{x} - \mathbf{x}_0|} + \frac{1}{|\mathbf{x} - \mathbf{x}_0|^2} P_i \frac{(\mathbf{x} - \mathbf{x}_0)_i}{|\mathbf{x} - \mathbf{x}_0|} + \frac{1}{|\mathbf{x} - \mathbf{x}_0|^3} Q_{ij} \left(\frac{3}{2} \frac{(\mathbf{x} - \mathbf{x}_0)_i (\mathbf{x} - \mathbf{x}_0)_j}{|\mathbf{x} - \mathbf{x}_0|} - \frac{1}{2} \delta_{ij} \right) + \dots \right] \quad (7.4)$$

where Einstein's summation convention over i and j is implied. As can be seen from this expression, only the first few terms are important if the field point (observation point) is far away from \mathbf{x}_0 .

For a normal medium, the major contributions to the electrostatic interactions come from the net charge and the lowest order electric multipole moments induced by the polarisation due to an applied electric field. Particularly important is the dipole moment. Let \mathbf{P} denote the electric dipole moment density (electric dipole moment per unit volume; unit: C/m²), also known as the *electric polarisation*, in some medium. In analogy with the second term in the expansion equation (7.4) on page 96, the electric potential from this volume distribution $\mathbf{P}(\mathbf{x}')$ of electric dipole moments \mathbf{p} at the source point \mathbf{x}' can be written

$$\begin{aligned} \phi_{\mathbf{p}}(\mathbf{x}) &= \frac{1}{4\pi\epsilon_0} \int_V \mathbf{P}(\mathbf{x}') \cdot \frac{\mathbf{x} - \mathbf{x}'}{|\mathbf{x} - \mathbf{x}'|^3} d^3x' \\ &= -\frac{1}{4\pi\epsilon_0} \int_V \mathbf{P}(\mathbf{x}') \cdot \nabla \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right) d^3x' \\ &= \frac{1}{4\pi\epsilon_0} \int_V \mathbf{P}(\mathbf{x}') \cdot \nabla' \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right) d^3x' \end{aligned} \quad (7.5)$$

Using the expression equation (M.68) on page 187 and applying the divergence theorem, we can rewrite this expression for the potential as follows:

$$\begin{aligned} \phi_{\mathbf{p}}(\mathbf{x}) &= \frac{1}{4\pi\epsilon_0} \left[\int_V \nabla' \cdot \left(\frac{\mathbf{P}(\mathbf{x}')}{|\mathbf{x} - \mathbf{x}'|} \right) d^3x' - \int_V \frac{\nabla' \cdot \mathbf{P}(\mathbf{x}')}{|\mathbf{x} - \mathbf{x}'|} d^3x' \right] \\ &= \frac{1}{4\pi\epsilon_0} \left[\oint_S \frac{\mathbf{P}(\mathbf{x}') \cdot \hat{\mathbf{n}}}{|\mathbf{x} - \mathbf{x}'|} d^2x - \int_V \frac{\nabla' \cdot \mathbf{P}(\mathbf{x}')}{|\mathbf{x} - \mathbf{x}'|} d^3x' \right] \end{aligned} \quad (7.6)$$

where the first term, which describes the effects of the induced, non-cancelling dipole moment on the surface of the volume, can be neglected, unless there is a discontinuity in $\mathbf{P} \cdot \hat{\mathbf{n}}$ at the surface. Doing so, we find that the contribution from

the electric dipole moments to the potential is given by

$$\phi_p = \frac{1}{4\pi\epsilon_0} \int_V \frac{-\nabla' \cdot \mathbf{P}(\mathbf{x}')}{|\mathbf{x} - \mathbf{x}'|} d^3x' \quad (7.7)$$

Comparing this expression with expression equation (3.3) on page 33 for the electrostatic potential from a static charge distribution ρ , we see that $-\nabla \cdot \mathbf{P}(\mathbf{x})$ has the characteristics of a charge density and that, to the lowest order, the effective charge density becomes $\rho(x) - \nabla \cdot \mathbf{P}(x)$, in which the second term is a polarisation term.

The version of equation equation (1.7) on page 4 where “true” and *polarisation charges* are separated thus becomes

$$\nabla \cdot \mathbf{E} = \frac{\rho(\mathbf{x}) - \nabla \cdot \mathbf{P}(\mathbf{x})}{\epsilon_0} \quad (7.8)$$

Rewriting this equation, and at the same time introducing the *electric displacement vector* (C/m^2)

$$\mathbf{D} = \epsilon_0 \mathbf{E} + \mathbf{P} \quad (7.9)$$

we obtain

$$\nabla \cdot (\epsilon_0 \mathbf{E} + \mathbf{P}) = \nabla \cdot \mathbf{D} = \rho^{\text{true}}(\mathbf{x}) \quad (7.10)$$

where ρ^{true} is the “true” charge density in the medium. This is one of Maxwell’s equations and is valid also for time varying fields. By introducing the notation $\rho^{\text{pol}} = -\nabla \cdot \mathbf{P}$ for the “polarised” charge density in the medium, and $\rho^{\text{total}} = \rho^{\text{true}} + \rho^{\text{pol}}$ for the “total” charge density, we can write down the following alternative version of Maxwell’s equation (7.23a) on page 100

$$\nabla \cdot \mathbf{E} = \frac{\rho^{\text{total}}(\mathbf{x})}{\epsilon_0} \quad (7.11)$$

Often, for low enough field strengths $|\mathbf{E}|$, the linear and isotropic relationship between \mathbf{P} and \mathbf{E}

$$\mathbf{P} = \epsilon_0 \chi \mathbf{E} \quad (7.12)$$

is a good approximation. The quantity χ is the *electric susceptibility* which is material dependent. For electromagnetically anisotropic media such as a magnetised plasma or a birefringent crystal, the susceptibility is a tensor. In general, the relationship is not of a simple linear form as in equation (7.12) but non-linear terms are important. In such a situation the principle of superposition is no longer valid and non-linear effects such as frequency conversion and mixing can be expected.

Inserting the approximation (7.12) into equation (7.9) above, we can write the

latter

$$\mathbf{D} = \epsilon \mathbf{E} \quad (7.13)$$

where, approximately,

$$\epsilon = \epsilon_0(1 + \chi) \quad (7.14)$$

7.2 Magnetisation and the magnetising field

An analysis of the properties of stationary magnetic media and the associated currents shows that three such types of currents exist:

1. In analogy with “true” charges for the electric case, we may have “true” currents \mathbf{j}^{true} , *i.e.*, a physical transport of true charges.
2. In analogy with electric polarisation \mathbf{P} there may be a form of charge transport associated with the changes of the polarisation with time. We call such currents induced by an external field *polarisation currents*. We identify them with $\partial \mathbf{P} / \partial t$.
3. There may also be intrinsic currents of a microscopic, often atomic, nature that are inaccessible to direct observation, but which may produce net effects at discontinuities and boundaries. We shall call such currents *magnetisation currents* and denote them \mathbf{j}_m .

No magnetic monopoles have been observed yet. So there is no correspondence in the magnetic case to the electric monopole moment (7.1). The lowest order magnetic moment, corresponding to the electric dipole moment (7.2), is the *magnetic dipole moment*

$$\mathbf{m} = \frac{1}{2} \int_V (\mathbf{x}' - \mathbf{x}_0) \times \mathbf{j}(\mathbf{x}') d^3x' \quad (7.15)$$

For a distribution of magnetic dipole moments in a volume, we may describe this volume in terms of the *magnetisation*, or magnetic dipole moment per unit volume, \mathbf{M} . Via the definition of the vector potential one can show that the magnetisation

current and the magnetisation is simply related:

$$\mathbf{j}_m = \nabla \times \mathbf{M} \quad (7.16)$$

In a stationary medium we therefore have a total current which is (approximately) the sum of the three currents enumerated above:

$$\mathbf{j}^{\text{total}} = \mathbf{j}^{\text{true}} + \frac{\partial \mathbf{P}}{\partial t} + \nabla \times \mathbf{M} \quad (7.17)$$

We then obtain the Maxwell equation

$$\nabla \times \mathbf{B} = \mu_0 \left(\mathbf{j}^{\text{true}} + \frac{\partial \mathbf{P}}{\partial t} + \nabla \times \mathbf{M} \right) \quad (7.18)$$

Moving the term $\nabla \times \mathbf{M}$ to the left hand side and introducing the *magnetising field* (*magnetic field intensity, Ampère-turn density*) as

$$\mathbf{H} = \frac{\mathbf{B}}{\mu_0} - \mathbf{M} \quad (7.19)$$

and using the definition for \mathbf{D} , equation (7.9) on page 97, we can write this Maxwell equation in the following form

$$\nabla \times \mathbf{H} = \mathbf{j}^{\text{true}} + \frac{\partial \mathbf{P}}{\partial t} = \mathbf{j}^{\text{true}} + \frac{\partial \mathbf{D}}{\partial t} - \epsilon_0 \frac{\partial \mathbf{E}}{\partial t} \quad (7.20)$$

We may, in analogy with the electric case, introduce a *magnetic susceptibility* for the medium. Denoting it χ_m , we can write

$$\mathbf{H} = \frac{\mathbf{B}}{\mu} \quad (7.21)$$

where, approximately,

$$\mu = \mu_0(1 + \chi_m) \quad (7.22)$$

7.3 Energy and momentum

As mentioned in chapter 1, Maxwell's equations expressed in terms of the derived field quantities \mathbf{D} and \mathbf{H} can be written

$$\nabla \cdot \mathbf{D} = \rho(t, \mathbf{x}) \quad (7.23a)$$

$$\nabla \cdot \mathbf{B} = 0 \quad (7.23b)$$

$$\nabla \times \mathbf{E} = -\frac{\partial \mathbf{B}}{\partial t} \quad (7.23c)$$

$$\nabla \times \mathbf{H} = \mathbf{j}(t, \mathbf{x}) + \frac{\partial \mathbf{D}}{\partial t} \quad (7.23d)$$

and are called *Maxwell's macroscopic equations*. These equations are convenient to use in certain simple cases. Together with the boundary conditions and the constitutive relations, they describe uniquely (but only approximately!) the properties of the electric and magnetic fields in matter. We shall use them in the following considerations on the energy and momentum of the electromagnetic field and its interaction with matter.

7.3.1 The energy theorem in Maxwell's theory

Scalar multiplying (7.23c) by \mathbf{H} , (7.23d) by \mathbf{E} and subtracting, we obtain

$$\begin{aligned} \mathbf{H} \cdot (\nabla \times \mathbf{E}) - \mathbf{E} \cdot (\nabla \times \mathbf{H}) &= \nabla \cdot (\mathbf{E} \times \mathbf{H}) \\ &= -\mathbf{H} \cdot \frac{\partial \mathbf{B}}{\partial t} - \mathbf{E} \cdot \mathbf{j} - \mathbf{E} \cdot \frac{\partial \mathbf{D}}{\partial t} = -\frac{1}{2} \frac{\partial}{\partial t} (\mathbf{H} \cdot \mathbf{B} + \mathbf{E} \cdot \mathbf{D}) - \mathbf{j} \cdot \mathbf{E} \end{aligned} \quad (7.24)$$

Integration over the entire volume V and using Gauss's theorem (the divergence theorem), we obtain

$$-\frac{\partial}{\partial t} \int_V \frac{1}{2} (\mathbf{H} \cdot \mathbf{B} + \mathbf{E} \cdot \mathbf{D}) d^3x' = \int_V \mathbf{j} \cdot \mathbf{E} d^3x' + \int_S (\mathbf{E} \times \mathbf{H}) \cdot d\mathbf{S}' \quad (7.25)$$

But, according to Ohm's law in the presence of an electromotive force field,

equation (1.26) on page 11:

$$\mathbf{j} = \sigma(\mathbf{E} + \mathbf{E}^{\text{EMF}}) \quad (7.26)$$

which means that

$$\int_V \mathbf{j} \cdot \mathbf{E} \, d^3x' = \int_V \frac{j^2}{\sigma} \, d^3x' - \int_V \mathbf{j} \cdot \mathbf{E}^{\text{EMF}} \, d^3x' \quad (7.27)$$

Inserting this into equation (7.25) on the facing page

$$\begin{aligned} \underbrace{\int_V \mathbf{j} \cdot \mathbf{E}^{\text{EMF}} \, d^3x'}_{\text{Applied electric power}} &= \underbrace{\int_V \frac{j^2}{\sigma} \, d^3x'}_{\text{Joule heat}} \\ &+ \underbrace{\frac{\partial}{\partial t} \int_V \frac{1}{2} (\mathbf{E} \cdot \mathbf{D} + \mathbf{H} \cdot \mathbf{B}) \, d^3x'}_{\text{Field energy}} \\ &+ \underbrace{\int_S (\mathbf{E} \times \mathbf{H}) \cdot d\mathbf{S}'}_{\text{Radiated power}} \end{aligned} \quad (7.28)$$

which is the *energy theorem in Maxwell's theory* also known as *Poynting's theorem*.

It is convenient to introduce the following quantities:

$$U_e = \frac{1}{2} \int_V \mathbf{E} \cdot \mathbf{D} \, d^3x' \quad (7.29)$$

$$U_m = \frac{1}{2} \int_V \mathbf{H} \cdot \mathbf{B} \, d^3x' \quad (7.30)$$

$$\mathbf{S} = \mathbf{E} \times \mathbf{H} \quad (7.31)$$

where U_e is the *electric field energy*, U_m is the *magnetic field energy*, both measured in J, and \mathbf{S} is the *Poynting vector (power flux)*, measured in W/m².

7.3.2 The momentum theorem in Maxwell's theory

Let us now investigate the momentum balance (force actions) in the case that a field interacts with matter in a non-relativistic way. For this purpose we consider the force density given by the *Lorentz force* per unit volume $\rho\mathbf{E} + \mathbf{j} \times \mathbf{B}$. Using

Maxwell's equations (7.23) and symmetrising, we obtain

$$\begin{aligned}
\rho \mathbf{E} + \mathbf{j} \times \mathbf{B} &= (\nabla \cdot \mathbf{D}) \mathbf{E} + \left(\nabla \times \mathbf{H} - \frac{\partial \mathbf{D}}{\partial t} \right) \times \mathbf{B} \\
&= \mathbf{E}(\nabla \cdot \mathbf{D}) + (\nabla \times \mathbf{H}) \times \mathbf{B} - \frac{\partial \mathbf{D}}{\partial t} \times \mathbf{B} \\
&= \mathbf{E}(\nabla \cdot \mathbf{D}) - \mathbf{B} \times (\nabla \times \mathbf{H}) \\
&\quad - \frac{\partial}{\partial t} (\mathbf{D} \times \mathbf{B}) + \mathbf{D} \times \frac{\partial \mathbf{B}}{\partial t} \\
&= \mathbf{E}(\nabla \cdot \mathbf{D}) - \mathbf{B} \times (\nabla \times \mathbf{H}) \\
&\quad - \frac{\partial}{\partial t} (\mathbf{D} \times \mathbf{B}) - \mathbf{D} \times (\nabla \times \mathbf{E}) + \mathbf{H}(\underbrace{\nabla \cdot \mathbf{B}}_{=0}) \\
&= [\mathbf{E}(\nabla \cdot \mathbf{D}) - \mathbf{D} \times (\nabla \times \mathbf{E})] + [\mathbf{H}(\nabla \cdot \mathbf{B}) - \mathbf{B} \times (\nabla \times \mathbf{H})] \\
&\quad - \frac{\partial}{\partial t} (\mathbf{D} \times \mathbf{B}) \tag{7.32}
\end{aligned}$$

One verifies easily that the i th vector components of the two terms in square brackets in the right hand member of (7.32) can be expressed as

$$\begin{aligned}
&[\mathbf{E}(\nabla \cdot \mathbf{D}) - \mathbf{D} \times (\nabla \times \mathbf{E})]_i \\
&= \frac{1}{2} \left(\mathbf{E} \cdot \frac{\partial \mathbf{D}}{\partial x_i} - \mathbf{D} \cdot \frac{\partial \mathbf{E}}{\partial x_i} \right) + \frac{\partial}{\partial x_j} \left(E_i D_j - \frac{1}{2} \mathbf{E} \cdot \mathbf{D} \delta_{ij} \right) \tag{7.33}
\end{aligned}$$

and

$$\begin{aligned}
&[\mathbf{H}(\nabla \cdot \mathbf{B}) - \mathbf{B} \times (\nabla \times \mathbf{H})]_i \\
&= \frac{1}{2} \left(\mathbf{H} \cdot \frac{\partial \mathbf{B}}{\partial x_i} - \mathbf{B} \cdot \frac{\partial \mathbf{H}}{\partial x_i} \right) + \frac{\partial}{\partial x_j} \left(H_i B_j - \frac{1}{2} \mathbf{B} \cdot \mathbf{H} \delta_{ij} \right) \tag{7.34}
\end{aligned}$$

respectively.

Using these two expressions in the i th component of equation (7.32) above and re-shuffling terms, we get

$$\begin{aligned}
(\rho \mathbf{E} + \mathbf{j} \times \mathbf{B})_i &- \frac{1}{2} \left[\left(\mathbf{E} \cdot \frac{\partial \mathbf{D}}{\partial x_i} - \mathbf{D} \cdot \frac{\partial \mathbf{E}}{\partial x_i} \right) \right. \\
&\quad \left. + \left(\mathbf{H} \cdot \frac{\partial \mathbf{B}}{\partial x_i} - \mathbf{B} \cdot \frac{\partial \mathbf{H}}{\partial x_i} \right) \right] + \frac{\partial}{\partial t} (\mathbf{D} \times \mathbf{B})_i \\
&= \frac{\partial}{\partial x_j} \left(E_i D_j - \frac{1}{2} \mathbf{E} \cdot \mathbf{D} \delta_{ij} + H_i B_j - \frac{1}{2} \mathbf{H} \cdot \mathbf{B} \delta_{ij} \right) \tag{7.35}
\end{aligned}$$

Introducing the *electric volume force* \mathbf{F}_{ev} via its i th component

$$(\mathbf{F}_{\text{ev}})_i = (\rho\mathbf{E} + \mathbf{j} \times \mathbf{B})_i - \frac{1}{2} \left[\left(\mathbf{E} \cdot \frac{\partial \mathbf{D}}{\partial x_i} - \mathbf{D} \cdot \frac{\partial \mathbf{E}}{\partial x_i} \right) + \left(\mathbf{H} \cdot \frac{\partial \mathbf{B}}{\partial x_i} - \mathbf{B} \cdot \frac{\partial \mathbf{H}}{\partial x_i} \right) \right] \quad (7.36)$$

and the *Maxwell stress tensor*

$$T_{ij} = E_i D_j - \frac{1}{2} \mathbf{E} \cdot \mathbf{D} \delta_{ij} + H_i B_j - \frac{1}{2} \mathbf{H} \cdot \mathbf{B} \delta_{ij} \quad (7.37)$$

we finally obtain the force equation

$$\left[\mathbf{F}_{\text{ev}} + \frac{\partial}{\partial t} (\mathbf{D} \times \mathbf{B}) \right]_i = \frac{\partial T_{ij}}{\partial x_j} \quad (7.38)$$

If we introduce the *relative electric permittivity* κ and the *relative magnetic permeability* κ_m as

$$\mathbf{D} = \kappa \epsilon_0 \mathbf{E} = \epsilon \mathbf{E} \quad (7.39)$$

$$\mathbf{B} = \kappa_m \mu_0 \mathbf{H} = \mu \mathbf{H} \quad (7.40)$$

we can rewrite (7.38) as

$$\frac{\partial T_{ij}}{\partial x_j} = \left(\mathbf{F}_{\text{ev}} + \frac{\kappa \kappa_m}{c^2} \frac{\partial \mathbf{S}}{\partial t} \right)_i \quad (7.41)$$

where \mathbf{S} is the Poynting vector defined in equation (7.31) on page 101. Integration over the entire volume V yields

$$\underbrace{\int_V \mathbf{F}_{\text{ev}} d^3x'}_{\text{Force on the matter}} + \frac{d}{dt} \underbrace{\int_V \frac{\kappa \kappa_m}{c^2} \mathbf{S} d^3x'}_{\text{Field momentum}} = \underbrace{\int_S \mathbf{T}_{\hat{n}} d^2x'}_{\text{Maxwell stress}} \quad (7.42)$$

which expresses the balance between the force on the matter, the rate of change of the electromagnetic field momentum and the Maxwell stress. This equation is called the *momentum theorem in Maxwell's theory*.

In vacuum (7.42) becomes

$$\int_V \rho(\mathbf{E} + \mathbf{v} \times \mathbf{B}) d^3x' + \frac{1}{c^2} \frac{d}{dt} \int_V \mathbf{S} d^3x' = \int_S \mathbf{T}_{\hat{n}} d^2x' \quad (7.43)$$

or

$$\frac{d}{dt} \mathbf{p}^{\text{mech}} + \frac{d}{dt} \mathbf{p}^{\text{field}} = \int_S \mathbf{T}_{\hat{n}} d^2x' \quad (7.44)$$

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Electromagnetic Radiation

In this chapter we will develop the theory of electromagnetic radiation, and therefore study electric and magnetic fields which are capable of carrying energy and momentum over large distances. In chapter 3 we were able to derive general expressions for the scalar and vector potentials from which we then, in chapter 4 calculated the total electric and magnetic fields from arbitrary distributions of charge and current sources. The only limitation in the calculation of the fields was that the advanced potentials were discarded.

We shall now study these fields further under the assumption that the observer is located in the *far zone*, *i.e.*, very far away from the source region(s). We therefore study the *radiation fields* which are dominating in this zone.

8.1 The radiation fields

From equation (4.13) on page 48 and equation (4.24) on page 51, which give the total electric and magnetic fields, we obtain the *radiation fields*

$$\begin{aligned} \mathbf{B}^{\text{rad}}(t, \mathbf{x}) &= \int_{-\infty}^{\infty} \mathbf{B}_{\omega}^{\text{rad}}(\mathbf{x}) e^{-i\omega t} d\omega \\ &= \frac{\mu_0}{4\pi c} \int_V \frac{\dot{\mathbf{j}}(t'_{\text{ret}}, \mathbf{x}') \times (\mathbf{x} - \mathbf{x}')}{|\mathbf{x} - \mathbf{x}'|^2} d^3x' \end{aligned} \quad (8.1)$$

$$\begin{aligned} \mathbf{E}^{\text{rad}}(t, \mathbf{x}) &= \int_{-\infty}^{\infty} \mathbf{E}_{\omega}^{\text{rad}}(\mathbf{x}) e^{-i\omega t} d\omega \\ &= \frac{1}{4\pi\epsilon_0 c^2} \int_V \frac{[\dot{\mathbf{j}}(t'_{\text{ret}}, \mathbf{x}') \times (\mathbf{x} - \mathbf{x}')]}{|\mathbf{x} - \mathbf{x}'|^3} \times (\mathbf{x} - \mathbf{x}') d^3x' \end{aligned} \quad (8.2)$$

where

$$\mathbf{j}(t'_{\text{ret}}, \mathbf{x}') \stackrel{\text{def}}{=} \left(\frac{\partial \mathbf{j}}{\partial t} \right)_{t=t'_{\text{ret}}} \quad (8.3)$$

Instead of studying the fields in the time domain, we can often make a spectrum analysis into the frequency domain and study each Fourier component separately. A superposition of all these components and a transformation back to the time domain will then yield the complete solution.

The Fourier representation of the radiation fields (8.1) (8.2) were included in equation (4.12) on page 48 and equation (4.23) on page 51, respectively and are explicitly given by

$$\begin{aligned} \mathbf{B}_\omega^{\text{rad}}(\mathbf{x}) &= \frac{1}{2\pi} \int_{-\infty}^{\infty} \mathbf{B}^{\text{rad}}(t, \mathbf{x}) e^{i\omega t} dt \\ &= -i \frac{k\mu_0}{4\pi} \int_V \frac{\mathbf{j}_\omega(\mathbf{x}') \times (\mathbf{x} - \mathbf{x}')}{|\mathbf{x} - \mathbf{x}'|^2} e^{ik|\mathbf{x} - \mathbf{x}'|} d^3x' \\ &= -i \frac{\mu_0}{4\pi} \int_V \frac{\mathbf{j}_\omega(\mathbf{x}') \times \mathbf{k}}{|\mathbf{x} - \mathbf{x}'|} e^{ik|\mathbf{x} - \mathbf{x}'|} d^3x' \end{aligned} \quad (8.4)$$

$$\begin{aligned} \mathbf{E}_\omega^{\text{rad}}(\mathbf{x}) &= \frac{1}{2\pi} \int_{-\infty}^{\infty} \mathbf{E}^{\text{rad}}(t, \mathbf{x}) e^{i\omega t} dt \\ &= -i \frac{k}{4\pi\epsilon_0 c} \int_V \frac{[\mathbf{j}_\omega(\mathbf{x}') \times (\mathbf{x} - \mathbf{x}')] \times (\mathbf{x} - \mathbf{x}')}{|\mathbf{x} - \mathbf{x}'|^3} e^{ik|\mathbf{x} - \mathbf{x}'|} d^3x' \\ &= -i \frac{1}{4\pi\epsilon_0 c} \int_V \frac{[\mathbf{j}_\omega(\mathbf{x}') \times \mathbf{k}] \times (\mathbf{x} - \mathbf{x}')}{|\mathbf{x} - \mathbf{x}'|^2} e^{ik|\mathbf{x} - \mathbf{x}'|} d^3x' \end{aligned} \quad (8.5)$$

If the source is located inside a volume V near \mathbf{x}_0 and has such a limited spatial extent that $\max |\mathbf{x}' - \mathbf{x}_0| \ll |\mathbf{x} - \mathbf{x}'|$, and the integration surface S , centred on \mathbf{x}_0 , has a large enough radius $|\mathbf{x} - \mathbf{x}_0| \gg \max |\mathbf{x}' - \mathbf{x}_0|$, we see from figure 8.1 on the next page that we can approximate

$$\begin{aligned} k|\mathbf{x} - \mathbf{x}'| &\equiv \mathbf{k} \cdot (\mathbf{x} - \mathbf{x}') \equiv \mathbf{k} \cdot (\mathbf{x} - \mathbf{x}_0) - \mathbf{k} \cdot (\mathbf{x}' - \mathbf{x}_0) \\ &\approx k|\mathbf{x} - \mathbf{x}_0| - \mathbf{k} \cdot (\mathbf{x}' - \mathbf{x}_0) \end{aligned} \quad (8.6)$$

If we recall from Formula (F.50) and formula (F.51) on page 166 that

$$dS = |\mathbf{x} - \mathbf{x}_0|^2 d\Omega = |\mathbf{x} - \mathbf{x}_0|^2 \sin \theta d\theta d\varphi$$

and note from figure 8.1 on the next page that $\hat{\mathbf{k}}$ and $\hat{\mathbf{n}}$ are nearly parallel, we see that we can approximate.

$$\frac{\hat{\mathbf{k}} \cdot d\mathbf{S}}{|\mathbf{x} - \mathbf{x}_0|^2} = \frac{\hat{\mathbf{k}} \cdot \hat{\mathbf{n}}}{|\mathbf{x} - \mathbf{x}_0|^2} dS \approx d\Omega \quad (8.7)$$

Both these approximations will be used in the following.

Within approximation (8.6) the expressions (8.4) and (8.5) for the radiation fields

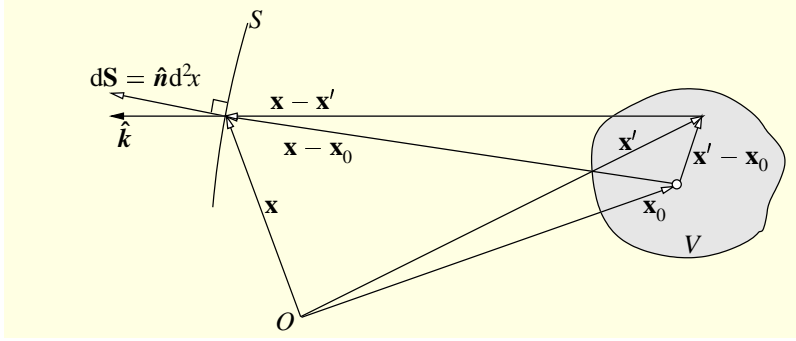


Figure 8.1. Relation between the surface normal and the \mathbf{k} vector for radiation generated at source points \mathbf{x}' near the point \mathbf{x}_0 in the source volume V . At distances much larger than the extent of V , the unit vector $\hat{\mathbf{n}}$, normal to the surface S which has its centre at \mathbf{x}_0 , and the unit vector $\hat{\mathbf{k}}$ of the radiation \mathbf{k} vector from \mathbf{x}' are nearly coincident.

can be approximated as

$$\begin{aligned} \mathbf{B}_\omega^{\text{rad}}(\mathbf{x}) &\approx -i \frac{\mu_0}{4\pi} e^{ik|\mathbf{x}-\mathbf{x}_0|} \int_V \frac{\mathbf{j}_\omega(\mathbf{x}') \times \mathbf{k}}{|\mathbf{x}-\mathbf{x}'|} e^{-ik \cdot (\mathbf{x}'-\mathbf{x}_0)} d^3x' \\ &\approx -i \frac{\mu_0}{4\pi} \frac{e^{ik|\mathbf{x}-\mathbf{x}_0|}}{|\mathbf{x}-\mathbf{x}_0|} \int_V [\mathbf{j}_\omega(\mathbf{x}') \times \mathbf{k}] e^{-ik \cdot (\mathbf{x}'-\mathbf{x}_0)} d^3x' \end{aligned} \quad (8.8)$$

$$\begin{aligned} \mathbf{E}_\omega^{\text{rad}}(\mathbf{x}) &\approx -i \frac{1}{4\pi\epsilon_0 c} e^{ik|\mathbf{x}-\mathbf{x}_0|} \int_V \frac{[\mathbf{j}_\omega(\mathbf{x}') \times \mathbf{k}] \times (\mathbf{x}-\mathbf{x}')}{|\mathbf{x}-\mathbf{x}'|^2} e^{-ik \cdot (\mathbf{x}'-\mathbf{x}_0)} d^3x' \\ &\approx i \frac{1}{4\pi\epsilon_0 c} \frac{e^{ik|\mathbf{x}-\mathbf{x}_0|}}{|\mathbf{x}-\mathbf{x}_0|} \frac{(\mathbf{x}-\mathbf{x}_0)}{|\mathbf{x}-\mathbf{x}_0|} \times \int_V [\mathbf{j}_\omega(\mathbf{x}') \times \mathbf{k}] e^{-ik \cdot (\mathbf{x}'-\mathbf{x}_0)} d^3x' \end{aligned} \quad (8.9)$$

I.e., if $\max |\mathbf{x}' - \mathbf{x}_0| \ll |\mathbf{x} - \mathbf{x}'|$, then the fields can be approximated as *spherical waves* multiplied by dimensional and angular factors, with integrals over points in the source volume only.

8.2 Radiated energy

Let us consider the energy that carried in the radiation fields \mathbf{B}^{rad} , Equation (8.1), and \mathbf{E}^{rad} , equation (8.2) on page 107. We have to treat signals with limited lifetime and hence finite frequency bandwidth differently from monochromatic signals.

8.2.1 Monochromatic signals

If the source is strictly monochromatic, we can obtain the temporal average of the radiated power P directly, simply by averaging over one period so that

$$\begin{aligned} \langle \mathbf{S} \rangle &= \langle \mathbf{E} \times \mathbf{H} \rangle = \frac{1}{2\mu_0} \text{Re} \{ \mathbf{E} \times \mathbf{B}^* \} = \frac{1}{2\mu_0} \text{Re} \{ \mathbf{E}_\omega e^{-i\omega t} \times (\mathbf{B}_\omega e^{-i\omega t})^* \} \\ &= \frac{1}{2\mu_0} \text{Re} \{ \mathbf{E}_\omega \times \mathbf{B}_\omega^* e^{-i\omega t} e^{i\omega t} \} = \frac{1}{2\mu_0} \text{Re} \{ \mathbf{E}_\omega \times \mathbf{B}_\omega^* \} \end{aligned} \quad (8.10)$$

Using the far-field approximations (8.8) and (8.9) and the fact that $1/c = \sqrt{\epsilon_0 \mu_0}$ and $R_0 = \sqrt{\mu_0/\epsilon_0}$ according to the definition (2.15) on page 25, we obtain

$$\langle \mathbf{S} \rangle = \frac{1}{32\pi^2} R_0 \frac{1}{|\mathbf{x} - \mathbf{x}_0|^2} \left| \int_V [(\mathbf{j}_\omega \times \mathbf{k})] e^{-i\mathbf{k} \cdot (\mathbf{x}' - \mathbf{x}_0)} d^3x' \right|^2 \frac{\mathbf{x} - \mathbf{x}_0}{|\mathbf{x} - \mathbf{x}_0|} \quad (8.11)$$

or, making use of (8.7) on page 108,

$$\frac{dP}{d\Omega} = \frac{1}{32\pi^2} R_0 \left| \int_V [(\mathbf{j}_\omega \times \mathbf{k})] e^{-i\mathbf{k} \cdot (\mathbf{x}' - \mathbf{x}_0)} d^3x' \right|^2 \quad (8.12)$$

which is the radiated power per unit solid angle.

8.2.2 Finite bandwidth signals

A signal with finite pulse width in time (t) domain has a certain spread in frequency (ω) domain. To calculate the total radiated energy we need to integrate over the whole bandwidth. The total energy transmitted through a unit area is the time integral of the Poynting vector:

$$\begin{aligned} \int_{-\infty}^{\infty} \mathbf{S}(t) dt &= \int_{-\infty}^{\infty} (\mathbf{E} \times \mathbf{H}) dt \\ &= \int_{-\infty}^{\infty} d\omega \int_{-\infty}^{\infty} d\omega' \int_{-\infty}^{\infty} (\mathbf{E}_\omega \times \mathbf{H}_{\omega'}) e^{-i(\omega+\omega')t} dt \end{aligned} \quad (8.13)$$

If we carry out the temporal integration first and use the fact that

$$\int_{-\infty}^{\infty} e^{-i(\omega+\omega')t} dt = 2\pi \delta(\omega + \omega') \quad (8.14)$$

equation (8.13) above can be written [*cf. Parseval's identity*]

$$\begin{aligned}
\int_{-\infty}^{\infty} \mathbf{S}(t) dt &= 2\pi \int_{-\infty}^{\infty} (\mathbf{E}_\omega \times \mathbf{H}_{-\omega}) d\omega \\
&= 2\pi \left(\int_0^{\infty} (\mathbf{E}_\omega \times \mathbf{H}_{-\omega}) d\omega + \int_{-\infty}^0 (\mathbf{E}_\omega \times \mathbf{H}_{-\omega}) d\omega \right) \\
&= 2\pi \left(\int_{\infty}^0 (\mathbf{E}_\omega \times \mathbf{H}_{-\omega}) d\omega - \int_0^{-\infty} (\mathbf{E}_\omega \times \mathbf{H}_{-\omega}) d\omega \right) \\
&= 2\pi \left(\int_0^{\infty} (\mathbf{E}_\omega \times \mathbf{H}_{-\omega}) d\omega + \int_0^{\infty} (\mathbf{E}_{-\omega} \times \mathbf{H}_\omega) d\omega \right) \\
&= \frac{2\pi}{\mu_0} \int_0^{\infty} (\mathbf{E}_\omega \times \mathbf{B}_{-\omega} + \mathbf{E}_{-\omega} \times \mathbf{B}_\omega) d\omega \\
&= \frac{2\pi}{\mu_0} \int_0^{\infty} (\mathbf{E}_\omega \times \mathbf{B}_\omega^* + \mathbf{E}_\omega^* \times \mathbf{B}_\omega) d\omega \tag{8.15}
\end{aligned}$$

where the last step follows from the real-valuedness of \mathbf{E}_ω and \mathbf{B}_ω . We insert the Fourier transforms of the field components which dominate at large distances, *i.e.*, the radiation fields (8.4) and (8.5).

The result, after integration over the area S of a large sphere which encloses the source, is

$$U = \frac{1}{4\pi} \sqrt{\frac{\mu_0}{\epsilon_0}} \int_S \int_0^{\infty} \left| \int_V \frac{\mathbf{j}_\omega \times \mathbf{k}}{|\mathbf{x} - \mathbf{x}'|} e^{ik|\mathbf{x} - \mathbf{x}'|} d^3x' \right|^2 \hat{\mathbf{k}} \cdot \hat{\mathbf{n}} dS d\omega \tag{8.16}$$

Inserting the approximations (8.6) and (8.7) into equation (8.16) and also introducing

$$U = \int_0^{\infty} U_\omega d\omega \tag{8.17}$$

and recalling the definition (2.15) on page 25 for the vacuum resistance R_0 we obtain

$$\frac{dU_\omega}{d\Omega} d\omega \approx \frac{1}{4\pi} R_0 \left| \int_V (\mathbf{j}_\omega \times \mathbf{k}) e^{-i\mathbf{k} \cdot (\mathbf{x}' - \mathbf{x}_0)} d^3x' \right|^2 d\omega \tag{8.18}$$

which, at large distances, is a good approximation to the energy that is radiated per unit solid angle $d\Omega$ in a frequency band $d\omega$. It is important to notice that Formula (8.18) includes only source coordinates. This means that the amount of energy that is being radiated is independent on the distance to the source (as long as it is large).

8.3 Radiation from extended sources

As shown above, one can, at least in principle, calculate the radiated fields, Poynting flux and energy for an arbitrary current density Fourier component. However, in practice, it is often difficult to evaluate the integrals unless the current has a simple distribution in space. In the general case, one has to resort to approximations. We shall consider both these situations.

Certain radiation systems have a geometry which is one-dimensional, symmetric or in any other way simple enough that a direct calculation of the radiated fields and energy is possible. This is for instance the case when the current flows in one direction in space only and is limited in extent. An example of this is a linear antenna.

8.3.1 Linear antenna

Let us apply equation (8.12) on page 110 for calculating the power from a linear, transmitting antenna, fed across a small gap at its centre with a monochromatic source. The antenna is a straight, thin conductor of length L which carries a one-dimensional time-varying current so that it produces electromagnetic radiation.

We assume that the conductor resistance and the energy loss due to the electromagnetic radiation are negligible. Since we can assume that the antenna wire is infinitely thin, the current must vanish at the end points $-L/2$ and $L/2$. The current therefore forms a standing wave with wave number $k = \omega/c$ and can be written

$$\mathbf{j}_0(\mathbf{x}') = I_0 \delta(x'_1) \delta(x'_2) \frac{\sin[k(L/2 - |x'_3|)]}{\sin(kL/2)} \hat{\mathbf{x}}_3 \quad (8.19)$$

where the amplitude I_0 is constant. In order to evaluate formula (8.12) on page 110 with the explicit monochromatic current (8.19) inserted, we need the expression

$$\begin{aligned} & \left| \int_V (\mathbf{j}_0 \times \mathbf{k}) e^{-i\mathbf{k} \cdot (\mathbf{x}' - \mathbf{x}_0)} d^3x' \right|^2 \\ &= \left| \int_{-L/2}^{L/2} I_0 \frac{\sin[k(L/2 - |x'_3|)]}{\sin(kL/2)} k \sin \theta e^{-ikx'_3 \cos \theta} e^{-ikx_0 \cos \theta} dx'_3 \right|^2 \\ &= I_0^2 \frac{k^2 \sin^2 \theta}{\sin^2(kL/2)} |e^{ikx_0 \cos \theta}|^2 \left| 2 \int_0^{L/2} \sin[k(L/2 - x'_3)] \cos(kx'_3 \cos \theta) dx'_3 \right|^2 \\ &= 4I_0^2 \left(\frac{\cos[(kL/2) \cos \theta] - \cos(kL/2)}{\sin \theta \sin(kL/2)} \right)^2 \quad (8.20) \end{aligned}$$

inserting this expression and $d\Omega = 2\pi \sin \theta d\theta$ into formula (8.12) on page 110

and integrating over θ , we find that the total radiated power from the antenna is

$$P(L) = R_0 I_0^2 \frac{1}{4\pi} \int_0^\pi \left(\frac{\cos[(kL/2) \cos \theta] - \cos(kL/2)}{\sin \theta \sin(kL/2)} \right)^2 \sin \theta \, d\theta \quad (8.21)$$

One can show that

$$\lim_{kL \rightarrow 0} P(L) = \frac{\pi}{12} \left(\frac{L}{\lambda} \right)^2 R_0 J_0^2 \quad (8.22)$$

where λ is the vacuum wavelength.

The quantity

$$R^{\text{rad}}(L) = \frac{P(L)}{I_{\text{eff}}^2} = \frac{P(L)}{\frac{1}{2} I_0^2} = R_0 \frac{\pi}{6} \left(\frac{L}{\lambda} \right)^2 \approx 197 \left(\frac{L}{\lambda} \right)^2 \Omega \quad (8.23)$$

is called the *radiation resistance*. For the technologically important case of a half-wave antenna, *i.e.*, for $L = \lambda/2$ or $kL = \pi$, formula (8.21) above reduces to

$$P(\lambda/2) = R_0 I_0^2 \frac{1}{4\pi} \int_0^\pi \frac{\cos^2(\frac{\pi}{2} \cos \theta)}{\sin \theta} \, d\theta \quad (8.24)$$

The integral in (8.24) can be evaluated numerically. It can also be evaluated analytically as follows:

$$\begin{aligned} \int_0^\pi \frac{\cos^2(\frac{\pi}{2} \cos \theta)}{\sin \theta} \, d\theta &= [\cos \theta \rightarrow u] = \int_{-1}^1 \frac{\cos^2(\frac{\pi}{2} v)}{1 - v^2} \, du = \\ &= \left[\cos^2\left(\frac{\pi}{2} v\right) = \frac{1 + \cos(\pi u)}{2} \right] \\ &= \frac{1}{2} \int_{-1}^1 \frac{1 + \cos(\pi u)}{(1+u)(1-u)} \, du \\ &= \frac{1}{4} \int_{-1}^1 \frac{1 + \cos(\pi u)}{(1+u)} \, du + \frac{1}{4} \int_{-1}^1 \frac{1 + \cos(\pi u)}{(1-u)} \, du \\ &= \frac{1}{2} \int_{-1}^1 \frac{1 + \cos(\pi u)}{(1+u)} \, du = \left[1 + u \rightarrow \frac{v}{\pi} \right] \\ &= \frac{1}{2} \int_0^{2\pi} \frac{1 - \cos v}{v} \, dv = \frac{1}{2} [\gamma + \ln 2\pi - \text{Ci}(2\pi)] \\ &\approx 1.22 \end{aligned} \quad (8.25)$$

where in the last step the *Euler-Mascheroni constant* $\gamma = 0.5772 \dots$ and the *cosine integral* $\text{Ci}(x)$ were introduced. Inserting this into the expression equation (8.24) we obtain the value $R^{\text{rad}}(\lambda/2) \approx 73 \Omega$.

8.4 Multipole radiation

In the general case, and when we are interested in evaluating the radiation far from the source volume, we can introduce an approximation which leads to a *multipole expansion* where individual terms can be evaluated analytically. We shall use *Hertz' method* to obtain this expansion.

8.4.1 The Hertz potential

Let us consider the continuity equation, which, according to expression (1.21) on page 9, can be written

$$\frac{\partial \rho(t, \mathbf{x})}{\partial t} + \nabla \cdot \mathbf{j}(t, \mathbf{x}) = 0 \quad (8.26)$$

If we introduce a vector field $\boldsymbol{\pi}(t, \mathbf{x})$ such that

$$\nabla \cdot \boldsymbol{\pi} = -\rho^{\text{true}} \quad (8.27a)$$

$$\frac{\partial \boldsymbol{\pi}}{\partial t} = \mathbf{j}^{\text{true}} \quad (8.27b)$$

and compare with equation (8.26), we see that $\boldsymbol{\pi}(t, \mathbf{x})$ satisfies this continuity equation. Furthermore, if we compare with the electric polarisation [*cf.* equation (7.9) on page 97], we see that the quantity $\boldsymbol{\pi}$ is related to the “true” charges in the same way as \mathbf{P} is related to polarised charge. Therefore, $\boldsymbol{\pi}$ is referred to as the *polarisation vector*.

We introduce a further potential \mathbf{II}^e with the following property

$$\nabla \cdot \mathbf{II}^e = -\phi \quad (8.28a)$$

$$\frac{1}{c^2} \frac{\partial \mathbf{II}^e}{\partial t} = \mathbf{A} \quad (8.28b)$$

where ϕ and \mathbf{A} are the electromagnetic scalar and vector potentials, respectively. As we see, \mathbf{II}^e acts as a “*super-potential*” in the sense that it is a potential from which we can obtain other potentials. It is called the *Hertz' vector* or *polarisation potential* and, as can be seen from (8.27) and (8.28), it satisfies the inhomogeneous wave equation

$$\square^2 \mathbf{II}^e = \frac{1}{c^2} \frac{\partial^2}{\partial t^2} \mathbf{II}^e - \nabla^2 \mathbf{II}^e = \frac{\boldsymbol{\pi}}{\epsilon_0} \quad (8.29)$$

This equation is of the same type as equation (3.19) on page 38, and has therefore

the retarded solution

$$\mathbf{\Pi}^e(t, \mathbf{x}) = \frac{1}{4\pi\epsilon_0} \int \frac{\boldsymbol{\pi}(t'_{\text{ret}}, \mathbf{x}')}{|\mathbf{x} - \mathbf{x}'|} d^3x' \quad (8.30)$$

with Fourier components

$$\mathbf{\Pi}^e_{\omega}(\mathbf{x}) = \frac{1}{4\pi\epsilon_0} \int \frac{\boldsymbol{\pi}_{\omega}(\mathbf{x}')e^{ik|\mathbf{x}-\mathbf{x}'|}}{|\mathbf{x} - \mathbf{x}'|} d^3x' \quad (8.31)$$

If we introduce the *help vector* \mathbf{C} such that

$$\mathbf{C} = \nabla \times \mathbf{\Pi}^e \quad (8.32)$$

we see that we can calculate the magnetic and electric fields, respectively, as follows

$$\mathbf{B} = \frac{1}{c^2} \frac{\partial \mathbf{C}}{\partial t} \quad (8.33a)$$

$$\mathbf{E} = \nabla \times \mathbf{C} \quad (8.33b)$$

Clearly, the last equation is valid only outside the source volume, where $\nabla \cdot \mathbf{E} = 0$. Since we are mainly interested in the fields in the far zone, a long distance from the source region, this is no essential limitation.

Assume that the source region is a limited volume around some central point \mathbf{x}_0 far away from the field (observation) point \mathbf{x} . Under these assumptions, we can expand expression (8.30) above the Hertz' vector, due to the presence of non-vanishing $\boldsymbol{\pi}(t'_{\text{ret}}, \mathbf{x}')$ in the vicinity of \mathbf{x}_0 , in a formal series. For this purpose we recall from *potential theory* that

$$\begin{aligned} \frac{e^{ik|\mathbf{x}-\mathbf{x}'|}}{|\mathbf{x} - \mathbf{x}'|} &\equiv \frac{e^{ik|(\mathbf{x}-\mathbf{x}_0)-(\mathbf{x}'-\mathbf{x}_0)|}}{|(\mathbf{x} - \mathbf{x}_0) - (\mathbf{x}' - \mathbf{x}_0)|} \\ &= ik \sum_{n=0}^{\infty} (2n+1)P_n(\cos \Theta) j_n(k|\mathbf{x}' - \mathbf{x}_0|) h_n^{(1)}(k|\mathbf{x} - \mathbf{x}_0|) \end{aligned} \quad (8.34)$$

where

$\frac{e^{ik|\mathbf{x}-\mathbf{x}'|}}{|\mathbf{x} - \mathbf{x}'|}$ is a *Green's function*

Θ is the angle between $\mathbf{x}' - \mathbf{x}_0$ and $\mathbf{x} - \mathbf{x}_0$

$P_n(\cos \Theta)$ is the *Legendre polynomial* of order n

$j_n(k|\mathbf{x}' - \mathbf{x}_0|)$ is the *spherical Bessel function of the first kind* of order n

$h_n^{(1)}(k|\mathbf{x} - \mathbf{x}_0|)$ is the *spherical Hankel function of the first kind* of order n

According to the addition theorem for Legendre polynomials, we can write

$$P_n(\cos \Theta) = \sum_{m=-n}^n (-1)^m P_n^m(\cos \theta) P_n^{-m}(\cos \theta') e^{im(\varphi - \varphi')} \quad (8.35)$$

where P_n^m is an *associated Legendre polynomial* and, in spherical polar coordinates,

$$\mathbf{x}' - \mathbf{x}_0 = (|\mathbf{x}' - \mathbf{x}_0|, \theta', \phi') \quad (8.36a)$$

$$\mathbf{x} - \mathbf{x}_0 = (|\mathbf{x} - \mathbf{x}_0|, \theta, \phi) \quad (8.36b)$$

Inserting equation (8.34) on the preceding page, together with equation (8.35) above, into equation (8.31) on the preceding page, we can in a formally exact way expand the Fourier component of the Hertz' vector as

$$\begin{aligned} \boldsymbol{\Pi}_\omega^e &= \frac{ik}{4\pi\epsilon_0} \sum_{n=0}^{\infty} \sum_{m=-n}^n (2n+1)(-1)^m h_n^{(1)}(k|\mathbf{x} - \mathbf{x}_0|) P_n^m(\cos \theta) e^{im\varphi} \\ &\quad \times \int_V \boldsymbol{\pi}_\omega(\mathbf{x}') j_n(k|\mathbf{x}' - \mathbf{x}_0|) P_n^{-m}(\cos \theta') e^{-im\varphi'} d^3x' \end{aligned} \quad (8.37)$$

We notice that there is no dependence on $\mathbf{x} - \mathbf{x}_0$ inside the integral; the integrand is only dependent on the relative source vector $\mathbf{x}' - \mathbf{x}_0$.

We are interested in the case where the field point is many wavelengths away from the well-localised sources, *i.e.*, when the following inequalities

$$k|\mathbf{x}' - \mathbf{x}_0| \ll 1 \ll k|\mathbf{x} - \mathbf{x}_0| \quad (8.38)$$

hold. Then we may to a good approximation replace $h_n^{(1)}$ with the first term in its asymptotic expansion:

$$h_n^{(1)}(k|\mathbf{x} - \mathbf{x}_0|) \approx (-i)^{n+1} \frac{e^{ik|\mathbf{x} - \mathbf{x}_0|}}{k|\mathbf{x} - \mathbf{x}_0|} \quad (8.39)$$

and replace j_n with the first term in its power series expansion:

$$j_n(k|\mathbf{x}' - \mathbf{x}_0|) \approx \frac{2^n n!}{(2n+1)!} (k|\mathbf{x}' - \mathbf{x}_0|)^n \quad (8.40)$$

Inserting these expansions into equation (8.37) above, we obtain the *multipole ex-*

ansion of the Fourier component of the Hertz' vector

$$\mathbf{\Pi}_\omega^e \approx \sum_{n=0}^{\infty} \mathbf{\Pi}_\omega^{e(n)} \quad (8.41a)$$

where

$$\mathbf{\Pi}_\omega^{e(n)} = (-i)^n \frac{1}{4\pi\epsilon_0} \frac{e^{ik|\mathbf{x}-\mathbf{x}_0|}}{|\mathbf{x}-\mathbf{x}_0|} \frac{2^n n!}{(2n)!} \int_V \boldsymbol{\pi}_\omega(\mathbf{x}') (k|\mathbf{x}'-\mathbf{x}_0|)^n P_n(\cos\Theta) d^3x' \quad (8.41b)$$

This expression is approximately correct only if certain care is exercised; if many $\mathbf{\Pi}_\omega^{e(n)}$ terms are needed for an accurate result, the expansions of the spherical Hankel and Bessel functions used above may not be consistent and must be replaced by more accurate expressions. Taking the inverse Fourier transform of $\mathbf{\Pi}_\omega^e$ will yield the Hertz' vector in time domain, which inserted into equation (8.32) on page 115 will yield \mathbf{C} . The resulting expression can then in turn be inserted into equation (8.33) on page 115 in order to obtain the radiation fields.

For a linear source distribution along the polar axis, $\Theta = \theta$ in expression (8.41b), and $P_n(\cos\theta)$ gives the angular distribution of the radiation. In the general case, however, the angular distribution must be computed with the help of formula (8.35) on the preceding page. Let us now study the lowest order contributions to the expansion of Hertz' vector.

8.4.2 Electric dipole radiation

Choosing $n = 0$ in expression (8.41b), we obtain

$$\mathbf{\Pi}_\omega^{e(0)} = \frac{e^{ik|\mathbf{x}-\mathbf{x}_0|}}{4\pi\epsilon_0 |\mathbf{x}-\mathbf{x}_0|} \int_V \boldsymbol{\pi}_\omega(\mathbf{x}') d^3x' = \frac{1}{4\pi\epsilon_0} \frac{e^{ik|\mathbf{x}-\mathbf{x}_0|}}{|\mathbf{x}-\mathbf{x}_0|} \mathbf{p}_\omega \quad (8.42)$$

where $\mathbf{p}_\omega = \int_V \boldsymbol{\pi}_\omega(\mathbf{x}') d^3x'$ is the Fourier component of the *electric dipole moment*; cf. equation (7.2) on page 95 which describes the static dipole moment. If a spherical coordinate system is chosen with its polar axis along \mathbf{p}_ω , the components of $\mathbf{\Pi}_\omega^{e(0)}$ are

$$\Pi_r = \Pi_\omega^{(0)} \cos\theta = \frac{1}{4\pi\epsilon_0} \frac{e^{ik|\mathbf{x}-\mathbf{x}_0|}}{|\mathbf{x}-\mathbf{x}_0|} p_\omega \cos\theta \quad (8.43a)$$

$$\Pi_\theta = -\Pi_\omega^{(0)} \sin\theta = -\frac{1}{4\pi\epsilon_0} \frac{e^{ik|\mathbf{x}-\mathbf{x}_0|}}{|\mathbf{x}-\mathbf{x}_0|} p_\omega \sin\theta \quad (8.43b)$$

$$\Pi_\phi = 0 \quad (8.43c)$$

Evaluating formula (8.32) on page 115 for the help vector \mathbf{C} , with the spherically

polar components (8.43) of $\mathbf{\Pi}_\omega^{e(0)}$ inserted, we obtain

$$\mathbf{C}_\omega = C_{\omega,\varphi}^{(0)} \hat{\boldsymbol{\phi}} = \frac{1}{4\pi\epsilon_0} \left(\frac{1}{|\mathbf{x} - \mathbf{x}_0|} - ik \right) \frac{e^{ik|\mathbf{x}-\mathbf{x}_0|}}{|\mathbf{x} - \mathbf{x}_0|} p_\omega \sin \theta \hat{\boldsymbol{\phi}} \quad (8.44)$$

Applying this to equation (8.33) on page 115, we obtain directly the Fourier components of the fields

$$\mathbf{B}_\omega = -i \frac{\omega\mu_0}{4\pi} \left(\frac{1}{|\mathbf{x} - \mathbf{x}_0|} - ik \right) \frac{e^{ik|\mathbf{x}-\mathbf{x}_0|}}{|\mathbf{x} - \mathbf{x}_0|} p_\omega \sin \theta \hat{\boldsymbol{\phi}} \quad (8.45)$$

$$\begin{aligned} \mathbf{E}_\omega = \frac{1}{4\pi\epsilon_0} \left[2 \left(\frac{1}{|\mathbf{x} - \mathbf{x}_0|^2} - \frac{ik}{|\mathbf{x} - \mathbf{x}_0|} \right) \cos \theta \frac{\mathbf{x} - \mathbf{x}_0}{|\mathbf{x} - \mathbf{x}_0|} \right. \\ \left. + \left(\frac{1}{|\mathbf{x} - \mathbf{x}_0|^2} - \frac{ik}{|\mathbf{x} - \mathbf{x}_0|} - k^2 \right) \sin \theta \hat{\boldsymbol{\theta}} \right] \frac{e^{ik|\mathbf{x}-\mathbf{x}_0|}}{|\mathbf{x} - \mathbf{x}_0|} p_\omega \end{aligned} \quad (8.46)$$

Keeping only those parts of the fields which dominate at large distances (the radiation fields) and recalling that the wave vector $\mathbf{k} = k(\mathbf{x} - \mathbf{x}_0)/|\mathbf{x} - \mathbf{x}_0|$ where $k = \omega/c$, we can now write down the Fourier components of the radiation parts of the magnetic and electric fields from the dipole:

$$\mathbf{B}_\omega^{\text{rad}} = -\frac{\omega\mu_0}{4\pi} \frac{e^{ik|\mathbf{x}-\mathbf{x}_0|}}{|\mathbf{x} - \mathbf{x}_0|} (\mathbf{p}_\omega \times \mathbf{k}) \quad (8.47a)$$

$$\mathbf{E}_\omega^{\text{rad}} = -\frac{1}{4\pi\epsilon_0} \frac{e^{ik|\mathbf{x}-\mathbf{x}_0|}}{|\mathbf{x} - \mathbf{x}_0|} [(\mathbf{p}_\omega \times \mathbf{k}) \times \mathbf{k}] \quad (8.47b)$$

These fields constitute the *electric dipole radiation*, also known as *E1 radiation*.

8.4.3 Magnetic dipole radiation

The next term in the expression (8.41b) on the preceding page for the expansion of the Fourier transform of the Hertz' vector is for $n = 1$:

$$\mathbf{\Pi}_\omega^{e(1)} = -i \frac{e^{ik|\mathbf{x}-\mathbf{x}_0|}}{4\pi\epsilon_0 |\mathbf{x} - \mathbf{x}_0|} \int_V k |\mathbf{x}' - \mathbf{x}_0| \boldsymbol{\pi}_\omega(\mathbf{x}') \cos \Theta d^3x' \quad (8.48)$$

$$= -ik \frac{1}{4\pi\epsilon_0 |\mathbf{x} - \mathbf{x}_0|^2} \int_V [(\mathbf{x} - \mathbf{x}_0) \cdot (\mathbf{x}' - \mathbf{x}_0)] \boldsymbol{\pi}_\omega(\mathbf{x}') d^3x' \quad (8.49)$$

Here, the term $[(\mathbf{x} - \mathbf{x}_0) \cdot (\mathbf{x}' - \mathbf{x}_0)] \boldsymbol{\pi}_\omega(\mathbf{x}')$ can be rewritten

$$[(\mathbf{x} - \mathbf{x}_0) \cdot (\mathbf{x}' - \mathbf{x}_0)] \boldsymbol{\pi}_\omega(\mathbf{x}') = (x_i - x_{0,i})(x'_i - x_{0,i}) \boldsymbol{\pi}_\omega(\mathbf{x}') \quad (8.50)$$

and introducing

$$\eta_i = x_i - x_{0,i} \quad (8.51a)$$

$$\eta'_i = x'_i - x_{0,i} \quad (8.51b)$$

the j th component of the integrand in $\mathbf{\Pi}_\omega^{e(1)}$ can be broken up into

$$\{[(\mathbf{x} - \mathbf{x}_0) \cdot (\mathbf{x}' - \mathbf{x}_0)] \boldsymbol{\pi}_\omega(\mathbf{x}')\}_j = \frac{1}{2} \eta_i (\pi_{\omega,j} \eta'_i + \pi_{\omega,i} \eta'_j) \quad (8.52a)$$

$$+ \frac{1}{2} \eta_i (\pi_{\omega,j} \eta'_i - \pi_{\omega,i} \eta'_j) \quad (8.52b)$$

i.e., as the sum of two parts, the first being symmetric and the second antisymmetric in the indices i, j . We note that the antisymmetric part can be written as

$$\begin{aligned} \frac{1}{2} \eta_i (\pi_{\omega,j} \eta'_i - \pi_{\omega,i} \eta'_j) &= \frac{1}{2} [\pi_{\omega,j} (\eta_i \eta'_i) - \eta'_j (\eta_i \pi_{\omega,j})] \\ &= \frac{1}{2} [\boldsymbol{\pi}_\omega (\boldsymbol{\eta} \cdot \boldsymbol{\eta}') - \boldsymbol{\eta}' (\boldsymbol{\eta} \cdot \boldsymbol{\pi}_\omega)]_j \\ &= \frac{1}{2} \{(\mathbf{x} - \mathbf{x}_0) \times [\boldsymbol{\pi}_\omega \times (\mathbf{x}' - \mathbf{x}_0)]\}_j \end{aligned} \quad (8.53)$$

equations (8.27) on page 114, and the fact that we are considering a single Fourier component,

$$\boldsymbol{\pi}(t, \mathbf{x}) = \boldsymbol{\pi}_\omega e^{-i\omega t} \quad (8.54)$$

allow us to express $\boldsymbol{\pi}_\omega$ in \mathbf{j}_ω as

$$\boldsymbol{\pi}_\omega = i \frac{\mathbf{j}_\omega}{\omega} \quad (8.55)$$

Hence, we can write the antisymmetric part of the integral in formula (8.49) on the facing page as

$$\begin{aligned} &\frac{1}{2} (\mathbf{x} - \mathbf{x}_0) \times \int_V \boldsymbol{\pi}_\omega(\mathbf{x}') \times (\mathbf{x}' - \mathbf{x}_0) dV' \\ &= i \frac{1}{2\omega} (\mathbf{x} - \mathbf{x}_0) \times \int_V \mathbf{j}_\omega(\mathbf{x}') \times (\mathbf{x}' - \mathbf{x}_0) d^3x' \\ &= -i \frac{1}{\omega} (\mathbf{x} - \mathbf{x}_0) \times \mathbf{m}_\omega \end{aligned} \quad (8.56)$$

where we introduced the Fourier transform of the *magnetic dipole moment*

$$\mathbf{m}_\omega = \frac{1}{2} \int_V (\mathbf{x}' - \mathbf{x}_0) \times \mathbf{j}_\omega(\mathbf{x}') d^3x' \quad (8.57)$$

The final result is that the antisymmetric, magnetic dipole, part of $\mathbf{\Pi}_\omega^{e(1)}$ can be

written

$$\mathbf{\Pi}_\omega^{\text{e,antisym}(1)} = -\frac{k}{4\pi\epsilon_0\omega} \frac{e^{ik|\mathbf{x}-\mathbf{x}_0|}}{|\mathbf{x}-\mathbf{x}_0|^2} (\mathbf{x}-\mathbf{x}_0) \times \mathbf{m}_\omega \quad (8.58)$$

In analogy with the electric dipole case, we insert this expression into equation (8.32) on page 115 to evaluate \mathbf{C} , with which Equations (8.33) on page 115 then gives the \mathbf{B} and \mathbf{E} fields. Discarding, as before, all terms belonging to the near fields and transition fields and keeping only the terms that dominate at large distances, we obtain

$$\mathbf{B}_\omega^{\text{rad}}(\mathbf{x}) = -\frac{\mu_0}{4\pi} \frac{e^{ik|\mathbf{x}-\mathbf{x}_0|}}{|\mathbf{x}-\mathbf{x}_0|} (\mathbf{m}_\omega \times \mathbf{k}) \times \mathbf{k} \quad (8.59a)$$

$$\mathbf{E}_\omega^{\text{rad}}(\mathbf{x}) = \frac{k}{4\pi\epsilon_0 c} \frac{e^{ik|\mathbf{x}-\mathbf{x}_0|}}{|\mathbf{x}-\mathbf{x}_0|} \mathbf{m}_\omega \times \mathbf{k} \quad (8.59b)$$

which are the fields of the *magnetic dipole radiation (M1 radiation)*.

8.4.4 Electric quadrupole radiation

The symmetric part $\mathbf{\Pi}_\omega^{\text{e,sym}(1)}$ of the $n = 1$ contribution in the equation (8.41b) on page 117 for the expansion of the Hertz' vector can be expressed in terms of the *electric quadrupole tensor*, which is defined in accordance with equation (7.3) on page 95:

$$\mathbf{Q}(t) = \int_V (\mathbf{x}' - \mathbf{x}_0)(\mathbf{x}' - \mathbf{x}_0)\rho(t, \mathbf{x}') d^3x' \quad (8.60)$$

Again we use this expression in equation (8.32) on page 115 to calculate the fields via Equations (8.33) on page 115. Tedious, but fairly straightforward algebra (which we will not present here), yields the resulting fields. The radiation components of the fields in the far field zone (wave zone) are given by

$$\mathbf{B}_\omega^{\text{rad}}(\mathbf{x}) = \frac{i\mu_0\omega}{8\pi} \frac{e^{ik|\mathbf{x}-\mathbf{x}_0|}}{|\mathbf{x}-\mathbf{x}_0|} (\mathbf{k} \cdot \mathbf{Q}_\omega) \times \mathbf{k} \quad (8.61a)$$

$$\mathbf{E}_\omega^{\text{rad}}(\mathbf{x}) = \frac{i}{8\pi\epsilon_0} \frac{e^{ik|\mathbf{x}-\mathbf{x}_0|}}{|\mathbf{x}-\mathbf{x}_0|} [(\mathbf{k} \cdot \mathbf{Q}_\omega) \times \mathbf{k}] \times \mathbf{k} \quad (8.61b)$$

This type of radiation is called *electric quadrupole radiation* or *E2 radiation*.

8.5 Radiation from a localised charge in arbitrary motion

The derivation of the radiation fields for the case of the source moving relative to the observer is considerably more complicated than the stationary cases studied above. In order to handle this non-stationary situation, we use the retarded potentials (3.36) on page 41 in chapter 3

$$\phi(t, \mathbf{x}) = \frac{1}{4\pi\epsilon_0} \int_V \frac{\rho(t'_{\text{ret}}, \mathbf{x}')}{|\mathbf{x} - \mathbf{x}'|} d^3x' \quad (8.62a)$$

$$\mathbf{A}(t, \mathbf{x}) = \frac{\mu_0}{4\pi} \int_V \frac{\mathbf{j}(t'_{\text{ret}}, \mathbf{x}')}{|\mathbf{x} - \mathbf{x}'|} d^3x' \quad (8.62b)$$

and consider a source region with such a limited spatial extent that the charges and currents are well localised. Specifically, we consider a charge q' , for instance an electron, which, classically, can be thought of as a localised, unstructured and rigid “charge distribution” with a small, finite radius. The part of this “charge distribution” dq' which we are considering is located in $dV' = d^3x'$ in the sphere in figure 8.5 on the following page. Since we assume that the electron (or any other similar electric charge) is moving with a velocity \mathbf{v} whose direction is arbitrary and whose magnitude can be almost comparable to the speed of light, we cannot say that the charge and current to be used in (8.62) is $\int_V \rho(t'_{\text{ret}}, \mathbf{x}') d^3x'$ and $\int_V \mathbf{v}\rho(t'_{\text{ret}}, \mathbf{x}') d^3x'$, respectively, because in the finite time interval during which the observed signal is generated, part of the charge distribution will “leak” out of the volume element d^3x' .

8.5.1 The Liénard-Wiechert potentials

The charge distribution in figure 8.5 on page 122 which contributes to the field at $\mathbf{x}(t)$ is located at $\mathbf{x}'(t')$ on a sphere with radius $r = |\mathbf{x} - \mathbf{x}'| = c(t - t')$. The radius interval of this sphere from which radiation is received at the field point \mathbf{x} during the time interval $(t, t + dt)$ is $(r, r + dr)$ and the net amount of charge in this radial interval is

$$dq' = \rho(t'_{\text{ret}}, \mathbf{x}') dS dr - \rho(t'_{\text{ret}}, \mathbf{x}') \frac{(\mathbf{x} - \mathbf{x}') \cdot \mathbf{v}}{|\mathbf{x} - \mathbf{x}'|} dS dt \quad (8.63)$$

where the last term represents the amount of “source leakage” due to the fact that the charge distribution moves with velocity \mathbf{v} . Since $dt = dr/c$ and $dS dr = d^3x'$

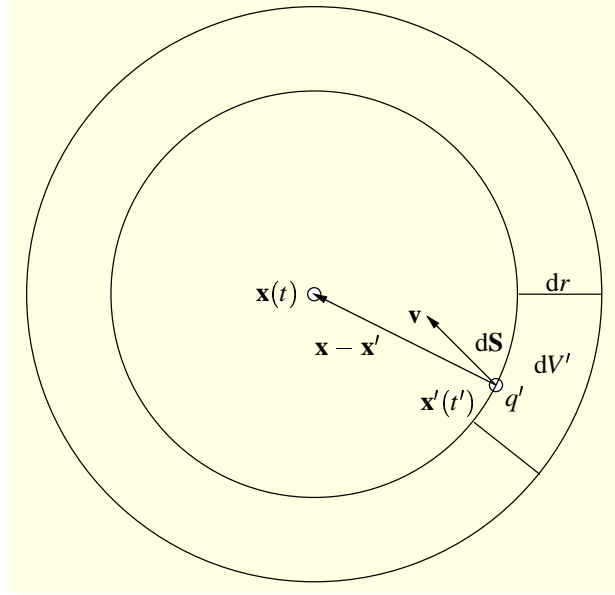


Figure 8.2. Signals which are observed at the field point \mathbf{x} at time t were generated at source points $\mathbf{x}'(t')$ on a sphere, centred on \mathbf{x} and expanding, as time increases, with the velocity \mathbf{c} outward from the centre. The source charge element moves with an arbitrary velocity \mathbf{v} and gives rise to a source “leakage” out of the source volume $dV' = d^3x'$.

we can rewrite this expression for the net charge as

$$\begin{aligned} dq' &= \rho(t'_{\text{ret}}, \mathbf{x}') d^3x' - \rho(t'_{\text{ret}}, \mathbf{x}') \frac{(\mathbf{x} - \mathbf{x}') \cdot \mathbf{v}}{c |\mathbf{x} - \mathbf{x}'|} d^3x' \\ &= \rho(t'_{\text{ret}}, \mathbf{x}') \left(1 - \frac{(\mathbf{x} - \mathbf{x}') \cdot \mathbf{v}}{c |\mathbf{x} - \mathbf{x}'|} \right) d^3x' \end{aligned} \quad (8.64)$$

or

$$\rho(t'_{\text{ret}}, \mathbf{x}') d^3x' = \frac{dq'}{1 - \frac{(\mathbf{x} - \mathbf{x}') \cdot \mathbf{v}}{c |\mathbf{x} - \mathbf{x}'|}} \quad (8.65)$$

which leads to the expression

$$\frac{\rho(t'_{\text{ret}}, \mathbf{x}')}{|\mathbf{x} - \mathbf{x}'|} d^3x' = \frac{dq'}{|\mathbf{x} - \mathbf{x}'| - \frac{(\mathbf{x} - \mathbf{x}') \cdot \mathbf{v}}{c}} \quad (8.66)$$

This is the expression to be used in the formulae (8.62) on the previous page for the retarded potentials. The result is (recall that $\mathbf{j} = \rho \mathbf{v}$)

$$\phi(t, \mathbf{x}) = \frac{1}{4\pi\epsilon_0} \int \frac{dq'}{|\mathbf{x} - \mathbf{x}'| - \frac{(\mathbf{x} - \mathbf{x}') \cdot \mathbf{v}}{c}} \quad (8.67a)$$

$$\mathbf{A}(t, \mathbf{x}) = \frac{\mu_0}{4\pi} \int \frac{\mathbf{v} dq'}{|\mathbf{x} - \mathbf{x}'| - \frac{(\mathbf{x} - \mathbf{x}') \cdot \mathbf{v}}{c}} \quad (8.67b)$$

For a sufficiently small and well localised charge distribution we can, assuming that the integrands do not change sign in the integration volume, use the mean value theorem and the fact that $\int_V dq' = q'$ to evaluate these expressions to become

$$\phi(t, \mathbf{x}) = \frac{q'}{4\pi\epsilon_0} \frac{1}{|\mathbf{x} - \mathbf{x}'| - \frac{(\mathbf{x} - \mathbf{x}') \cdot \mathbf{v}}{c}} = \frac{q'}{4\pi\epsilon_0} \frac{1}{s} \quad (8.68a)$$

$$\mathbf{A}(t, \mathbf{x}) = \frac{q'}{4\pi\epsilon_0 c^2} \frac{\mathbf{v}}{|\mathbf{x} - \mathbf{x}'| - \frac{(\mathbf{x} - \mathbf{x}') \cdot \mathbf{v}}{c}} = \frac{q'}{4\pi\epsilon_0 c^2} \frac{\mathbf{v}}{s} = \frac{\mathbf{v}}{c^2} \phi(t, \mathbf{x}) \quad (8.68b)$$

where

$$s = |\mathbf{x} - \mathbf{x}'| - \frac{(\mathbf{x} - \mathbf{x}') \cdot \mathbf{v}}{c} \quad (8.69a)$$

$$= |\mathbf{x} - \mathbf{x}'| \left(1 - \frac{\mathbf{x} - \mathbf{x}'}{|\mathbf{x} - \mathbf{x}'|} \cdot \frac{\mathbf{v}}{c} \right) \quad (8.69b)$$

$$= (\mathbf{x} - \mathbf{x}') \cdot \left(\frac{\mathbf{x} - \mathbf{x}'}{|\mathbf{x} - \mathbf{x}'|} - \frac{\mathbf{v}}{c} \right) \quad (8.69c)$$

is the *retarded relative distance*. The potentials (8.68) are precisely the *Liénard-Wiechert potentials* which we derived in section 5.3.2 on page 67 by using a covariant formalism.

It is important to realise that in the complicated derivation presented here, the observer is in a coordinate system which has an “absolute” meaning and the velocity \mathbf{v} is that of the particle, whereas in the covariant derivation two frames of equal standing were moving relative to each other with \mathbf{v} . Expressed in the four-potential, equation (5.46) on page 66, the Liénard-Wiechert potentials become

$$A^\mu(x^\kappa) = \frac{q'}{4\pi\epsilon_0} \left(\frac{1}{s}, \frac{\mathbf{v}}{cs} \right) = (\phi, c\mathbf{A}) \quad (8.70)$$

The Liénard-Wiechert potentials are applicable to all problems where a spatially localised charge emits electromagnetic radiation, and we shall now study such emission problems. The electric and magnetic fields are calculated from the

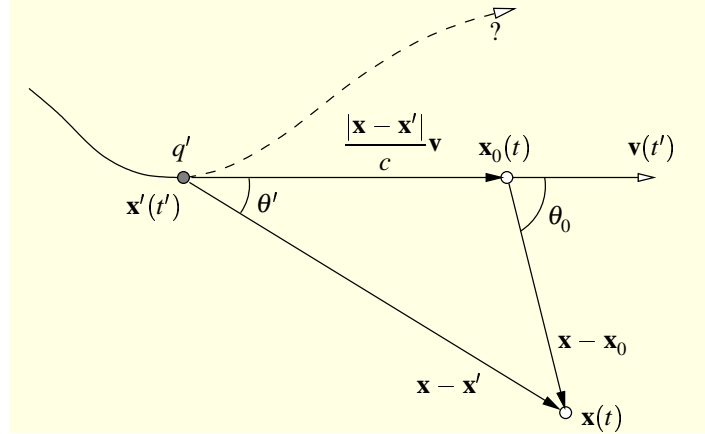


Figure 8.3. Signals which are observed at the field point \mathbf{x} at time t were generated at the source point $\mathbf{x}'(t')$. After time t' the particle, which moves with nonuniform velocity, has followed a yet unknown trajectory. Extrapolating tangentially the trajectory from $\mathbf{x}'(t')$, based on the velocity $\mathbf{v}(t')$, defines the *virtual simultaneous coordinate* $\mathbf{x}_0(t)$.

potentials in the usual way:

$$\mathbf{B}(t, \mathbf{x}) = \nabla \times \mathbf{A}(t, \mathbf{x}) \quad (8.71a)$$

$$\mathbf{E}(t, \mathbf{x}) = -\nabla\phi(t, \mathbf{x}) - \frac{\partial \mathbf{A}(t, \mathbf{x})}{\partial t} \quad (8.71b)$$

8.5.2 Radiation from an accelerated point charge

Consider a localised charge q' and assume that its trajectory is known experimentally as a function of *retarded time*

$$\mathbf{x}' = \mathbf{x}'(t') \quad (8.72)$$

(in the interest of simplifying our notation, we drop the subscript “ret” on t' from now on). This means that we know the trajectory of the charge q' , *i.e.*, \mathbf{x}' , for all times up to the time t' at which a signal was emitted in order to precisely arrive at the field point \mathbf{x} at time t . Because of the finite speed of propagation of the fields, the trajectory at times later than t' is not (yet) known.

The retarded velocity and acceleration at time t' are given by

$$\mathbf{v}(t') = \frac{d\mathbf{x}'}{dt'} \quad (8.73a)$$

$$\mathbf{a}(t') = \dot{\mathbf{v}}(t') = \frac{d\mathbf{v}}{dt'} = \frac{d^2\mathbf{x}'}{dt'^2} \quad (8.73b)$$

As for the charge coordinate \mathbf{x}' itself, we have in general no knowledge of the velocity and acceleration at times later than t' , in particular not at the time of observation t . If we choose the field point \mathbf{x} as fixed, application of (8.73) to the relative vector $\mathbf{x} - \mathbf{x}'$ yields

$$\frac{d}{dt'}(\mathbf{x} - \mathbf{x}'(t')) = -\mathbf{v}(t') \quad (8.74a)$$

$$\frac{d^2}{dt'^2}(\mathbf{x} - \mathbf{x}'(t')) = -\dot{\mathbf{v}}(t') \quad (8.74b)$$

The retarded time t' can, at least in principle, be calculated from the implicit relation

$$t' = t'(t, \mathbf{x}) = t - \frac{|\mathbf{x} - \mathbf{x}'(t')|}{c} \quad (8.75)$$

and we shall see later how this relation can be taken into account in the calculations.

According to formulae (8.71) on the facing page the electric and magnetic fields are determined via differentiation of the retarded potentials at the observation time t and at the observation point \mathbf{x} . In these formulae the unprimed ∇ , *i.e.*, the spatial derivative differentiation operator $\nabla = \hat{x}_i \partial / \partial x_i$ means that we differentiate with respect to the coordinates $\mathbf{x} = (x_1, x_2, x_3)$ while keeping t fixed, and the unprimed time derivative operator $\partial / \partial t$ means that we differentiate with respect to t while keeping \mathbf{x} fixed. But the Liénard-Wiechert potentials ϕ and \mathbf{A} , equations (8.68) on page 123, are expressed in the charge velocity $\mathbf{v}(t')$ given by equation (8.73a) and the retarded relative distance $s(t', \mathbf{x})$ given by equation (8.69) on page 123. This means that the expressions for the potentials ϕ and \mathbf{A} contain terms which are expressed explicitly in t' , which in turn is expressed implicitly in t via equation (8.75). Despite this complication it is possible, as we shall see below, to determine the electric and magnetic fields and associated quantities at the time of observation t . To this end, we need to investigate carefully the action of differentiation on the potentials.

The differential operator method We introduce the convention that a differential operator embraced by parentheses with an index \mathbf{x} or t means that the operator in question is applied at constant \mathbf{x} and t , respectively. With this convention, we

find that

$$\left(\frac{\partial}{\partial t'}\right)_{\mathbf{x}} |\mathbf{x} - \mathbf{x}'(t')| = \frac{\mathbf{x} - \mathbf{x}'}{|\mathbf{x} - \mathbf{x}'|} \cdot \left(\frac{\partial}{\partial t'}\right)_{\mathbf{x}} (\mathbf{x} - \mathbf{x}'(t')) = -\frac{(\mathbf{x} - \mathbf{x}') \cdot \mathbf{v}}{|\mathbf{x} - \mathbf{x}'|} \quad (8.76)$$

Furthermore, by applying the operator $(\partial/\partial t)_{\mathbf{x}}$ to equation (8.75) on the previous page we find that

$$\begin{aligned} \left(\frac{\partial t'}{\partial t}\right)_{\mathbf{x}} &= 1 - \left(\frac{\partial}{\partial t}\right)_{\mathbf{x}} \frac{|\mathbf{x} - \mathbf{x}'(t'(t, \mathbf{x}))|}{c} \\ &= 1 - \left[\left(\frac{\partial}{\partial t'}\right)_{\mathbf{x}} \frac{|\mathbf{x} - \mathbf{x}'|}{c}\right] \left(\frac{\partial t'}{\partial t}\right)_{\mathbf{x}} \\ &= 1 + \frac{(\mathbf{x} - \mathbf{x}') \cdot \mathbf{v}}{c|\mathbf{x} - \mathbf{x}'|} \left(\frac{\partial t'}{\partial t}\right)_{\mathbf{x}} \end{aligned} \quad (8.77)$$

This is an algebraic equation in $(\partial t'/\partial t)_{\mathbf{x}}$ which we can solve to obtain

$$\left(\frac{\partial t'}{\partial t}\right)_{\mathbf{x}} = \frac{|\mathbf{x} - \mathbf{x}'|}{|\mathbf{x} - \mathbf{x}'| - (\mathbf{x} - \mathbf{x}') \cdot \mathbf{v}/c} = \frac{|\mathbf{x} - \mathbf{x}'|}{s} \quad (8.78)$$

where $s = s(t', \mathbf{x})$ is the retarded relative distance given by equation (8.69) on page 123. Making use of equation (8.78) above, we obtain the following useful operator identity

$$\left(\frac{\partial}{\partial t}\right)_{\mathbf{x}} = \left(\frac{\partial t'}{\partial t}\right)_{\mathbf{x}} \left(\frac{\partial}{\partial t'}\right)_{\mathbf{x}} = \frac{|\mathbf{x} - \mathbf{x}'|}{s} \left(\frac{\partial}{\partial t'}\right)_{\mathbf{x}} \quad (8.79)$$

Likewise, by applying $(\nabla)_{\mathbf{t}}$ to equation (8.75) on the preceding page we obtain

$$\begin{aligned} (\nabla)_{\mathbf{t}} t' &= -(\nabla)_{\mathbf{t}} \frac{|\mathbf{x} - \mathbf{x}'(t'(t, \mathbf{x}))|}{c} = -\frac{\mathbf{x} - \mathbf{x}'}{c|\mathbf{x} - \mathbf{x}'|} \cdot (\nabla)_{\mathbf{t}} (\mathbf{x} - \mathbf{x}') \\ &= -\frac{\mathbf{x} - \mathbf{x}'}{c|\mathbf{x} - \mathbf{x}'|} + \frac{(\mathbf{x} - \mathbf{x}') \cdot \mathbf{v}}{c|\mathbf{x} - \mathbf{x}'|} (\nabla)_{\mathbf{t}} t' \end{aligned} \quad (8.80)$$

This is an algebraic equation in $(\nabla)_{\mathbf{t}}$ with the solution

$$(\nabla)_{\mathbf{t}} t' = -\frac{\mathbf{x} - \mathbf{x}'}{cs} \quad (8.81)$$

which gives the following operator relation when $(\nabla)_{\mathbf{t}}$ is acting on an arbitrary function of t' and \mathbf{x} :

$$(\nabla)_{\mathbf{t}} = [(\nabla)_{\mathbf{t}} t'] \left(\frac{\partial}{\partial t'}\right)_{\mathbf{x}} + (\nabla)_{\mathbf{t}'} = -\frac{\mathbf{x} - \mathbf{x}'}{cs} \left(\frac{\partial}{\partial t'}\right)_{\mathbf{x}} + (\nabla)_{\mathbf{t}'} \quad (8.82)$$

With the help of the rules (8.82) and (8.79) we are now able to replace t by t' in the

operations which we need to perform. We find, for instance, that

$$\begin{aligned}\nabla\phi &\equiv (\nabla\phi)_t = \nabla\left(\frac{1}{4\pi\epsilon_0}\frac{q'}{s}\right) \\ &= -\frac{q'}{4\pi\epsilon_0 s^2}\left[\frac{\mathbf{x}-\mathbf{x}'}{|\mathbf{x}-\mathbf{x}'|} - \frac{\mathbf{v}}{c} - \frac{\mathbf{x}-\mathbf{x}'}{cs}\left(\frac{\partial s}{\partial t'}\right)_\mathbf{x}\right]\end{aligned}\quad (8.83a)$$

$$\begin{aligned}\frac{\partial\mathbf{A}}{\partial t} &\equiv \left(\frac{\partial\mathbf{A}}{\partial t}\right)_\mathbf{x} = \frac{\partial}{\partial t}\left(\frac{\mu_0}{4\pi}\frac{q'\mathbf{v}}{s}\right) \\ &= \frac{q'}{4\pi\epsilon_0 c^2 s^3}\left[|\mathbf{x}-\mathbf{x}'|s\dot{\mathbf{v}} - |\mathbf{x}-\mathbf{x}'|\mathbf{v}\left(\frac{\partial s}{\partial t'}\right)_\mathbf{x}\right]\end{aligned}\quad (8.83b)$$

Utilising these relations in the calculation of the \mathbf{E} field from the Liénard-Wiechert potentials, equations (8.68) on page 123, we obtain

$$\begin{aligned}\mathbf{E}(t, \mathbf{x}) &= -\nabla\phi - \frac{\partial\mathbf{A}}{\partial t} \\ &= \frac{q'}{4\pi\epsilon_0 s^2}\left[\frac{(\mathbf{x}-\mathbf{x}') - |\mathbf{x}-\mathbf{x}'|\mathbf{v}/c}{|\mathbf{x}-\mathbf{x}'|} - \frac{(\mathbf{x}-\mathbf{x}') - |\mathbf{x}-\mathbf{x}'|\mathbf{v}/c}{cs}\left(\frac{\partial s}{\partial t'}\right)_\mathbf{x} - \frac{|\mathbf{x}-\mathbf{x}'|\dot{\mathbf{v}}}{c^2}\right]\end{aligned}\quad (8.84)$$

Starting from expression (8.69a) on page 123 for the retarded relative distance $s(t', \mathbf{x})$, we see that we can evaluate $(\partial s/\partial t')_\mathbf{x}$ in the following way

$$\begin{aligned}\left(\frac{\partial s}{\partial t'}\right)_\mathbf{x} &= \left(\frac{\partial}{\partial t'}\right)_\mathbf{x}\left(|\mathbf{x}-\mathbf{x}'| - \frac{(\mathbf{x}-\mathbf{x}')\cdot\mathbf{v}}{c}\right) \\ &= \frac{\partial}{\partial t'}|\mathbf{x}-\mathbf{x}'(t')| - \frac{1}{c}\left(\frac{\partial(\mathbf{x}-\mathbf{x}'(t'))}{\partial t'}\cdot\mathbf{v} - (\mathbf{x}-\mathbf{x}')\cdot\frac{\partial\mathbf{v}(t')}{\partial t'}\right) \\ &= -\frac{(\mathbf{x}-\mathbf{x}')\cdot\mathbf{v}}{|\mathbf{x}-\mathbf{x}'|} + \frac{v^2}{c} - \frac{(\mathbf{x}-\mathbf{x}')\cdot\dot{\mathbf{v}}}{c}\end{aligned}\quad (8.85)$$

where equation (8.76) on the preceding page and equations (8.73) on page 124, respectively, were used. Hence, the electric field generated by an arbitrarily moving charged particle at $\mathbf{x}'(t')$ is given by the expression

$$\mathbf{E}(t, \mathbf{x}) = \frac{q'}{4\pi\epsilon_0 s^3} \left\{ \underbrace{\left((\mathbf{x} - \mathbf{x}') - \frac{|\mathbf{x} - \mathbf{x}'| \mathbf{v}}{c} \right) \left(1 - \frac{v^2}{c^2} \right)}_{\text{Coulomb field when } v \rightarrow 0} + \underbrace{\frac{\mathbf{x} - \mathbf{x}'}{c^2} \times \left[\left((\mathbf{x} - \mathbf{x}') - \frac{|\mathbf{x} - \mathbf{x}'| \mathbf{v}}{c} \right) \times \dot{\mathbf{v}} \right]}_{\text{Radiation field}} \right\} \quad (8.86)$$

The first part of the field, the *velocity field*, tends to the ordinary Coulomb field when $v \rightarrow 0$ and does not contribute to the radiation. The second part of the field, the *acceleration field*, is radiated into the far zone and is therefore also called the *radiation field*.

From Figure 8.3 we see that the position the charged particle would have had if at t' all external forces would have been switched off so that the trajectory from then on would have been a straight line in the direction of the tangent at $\mathbf{x}'(t')$ is $\mathbf{x}_0(t)$, the *virtual simultaneous coordinate*. During the arbitrary motion, we interpret $\mathbf{x} - \mathbf{x}_0$ as the coordinate of the field point \mathbf{x} relative to the virtual simultaneous coordinate $\mathbf{x}_0(t)$. Since the time it takes from a signal to propagate (in the assumed vacuum) from $\mathbf{x}'(t')$ to \mathbf{x} is $|\mathbf{x} - \mathbf{x}'|/c$, this relative vector is given by

$$\mathbf{x} - \mathbf{x}_0 = (\mathbf{x} - \mathbf{x}') - \frac{|\mathbf{x} - \mathbf{x}'| \mathbf{v}}{c} \quad (8.87)$$

This allows us to rewrite equation (8.86) above in the following way

$$\mathbf{E}(t, \mathbf{x}) = \frac{q'}{4\pi\epsilon_0 s^3} \left[(\mathbf{x} - \mathbf{x}_0) \left(1 - \frac{v^2}{c^2} \right) + (\mathbf{x} - \mathbf{x}') \times \frac{(\mathbf{x} - \mathbf{x}_0) \times \dot{\mathbf{v}}}{c^2} \right] \quad (8.88)$$

In a similar manner we can compute the magnetic field:

$$\begin{aligned} \mathbf{B}(t, \mathbf{x}) &= \nabla \times \mathbf{A} \equiv (\nabla)_t \times \mathbf{A} = (\nabla)_{t'} \times \mathbf{A} - \frac{\mathbf{x} - \mathbf{x}'}{cs} \times \left(\frac{\partial}{\partial t'} \right)_x \mathbf{A} \\ &= -\frac{q'}{4\pi\epsilon_0 c^2 s^2} \frac{\mathbf{x} - \mathbf{x}'}{|\mathbf{x} - \mathbf{x}'|} \times \mathbf{v} - \frac{\mathbf{x} - \mathbf{x}'}{c|\mathbf{x} - \mathbf{x}'|} \times \left(\frac{\partial \mathbf{A}}{\partial t} \right)_x \end{aligned} \quad (8.89)$$

where we made use of equation (8.68) on page 123 and formula (8.79) on page 126. But, according to (8.83a),

$$\frac{\mathbf{x} - \mathbf{x}'}{c|\mathbf{x} - \mathbf{x}'|} \times (\nabla)_t \phi = \frac{q'}{4\pi\epsilon_0 c^2 s^2} \frac{\mathbf{x} - \mathbf{x}'}{|\mathbf{x} - \mathbf{x}'|} \times \mathbf{v} \quad (8.90)$$

so that

$$\begin{aligned}
\mathbf{B}(t, \mathbf{x}) &= \frac{\mathbf{x} - \mathbf{x}'}{c|\mathbf{x} - \mathbf{x}'|} \times \left[-(\nabla\phi)_t - \left(\frac{\partial \mathbf{A}}{\partial t} \right)_{\mathbf{x}} \right] \\
&= \frac{\mathbf{x} - \mathbf{x}'}{c|\mathbf{x} - \mathbf{x}'|} \times \mathbf{E}(t, \mathbf{x})
\end{aligned} \tag{8.91}$$

The radiation part of the electric field is obtained from the acceleration field in formula (8.86) on the facing page as

$$\begin{aligned}
\mathbf{E}^{\text{rad}}(t, \mathbf{x}) &= \lim_{|\mathbf{x} - \mathbf{x}'| \rightarrow \infty} \mathbf{E}(t, \mathbf{x}) \\
&= \frac{q'}{4\pi\epsilon_0 c^2 s^3} (\mathbf{x} - \mathbf{x}') \times \left[\left((\mathbf{x} - \mathbf{x}') - \frac{|\mathbf{x} - \mathbf{x}'| \mathbf{v}}{c} \right) \times \dot{\mathbf{v}} \right] \\
&= \frac{q'}{4\pi\epsilon_0 c^2 s^3} (\mathbf{x} - \mathbf{x}') \times [(\mathbf{x} - \mathbf{x}_0) \times \dot{\mathbf{v}}]
\end{aligned} \tag{8.92}$$

where in the last step we again used formula (8.87) on the preceding page. Using this formula and formula (8.91), the radiation part of the magnetic field can be written

$$\mathbf{B}^{\text{rad}}(t, \mathbf{x}) = \frac{\mathbf{x} - \mathbf{x}'}{c|\mathbf{x} - \mathbf{x}'|} \times \mathbf{E}^{\text{rad}}(t, \mathbf{x}) \tag{8.93}$$

The direct method An alternative to the differential operator transformation technique just described is to try to express all quantities in the potentials directly in t and \mathbf{x} . An example of such a quantity is the retarded relative distance $s(t', \mathbf{x})$. According to equation (8.69) on page 123, the square of this retarded relative can be written

$$s^2(t', \mathbf{x}) = |\mathbf{x} - \mathbf{x}'|^2 - 2|\mathbf{x} - \mathbf{x}'| \frac{(\mathbf{x} - \mathbf{x}') \cdot \mathbf{v}}{c} + \left(\frac{(\mathbf{x} - \mathbf{x}') \cdot \mathbf{v}}{c} \right)^2 \tag{8.94}$$

If we use the following handy identity

$$\begin{aligned}
&\left(\frac{(\mathbf{x} - \mathbf{x}') \cdot \mathbf{v}}{c} \right)^2 + \left(\frac{(\mathbf{x} - \mathbf{x}') \times \mathbf{v}}{c} \right)^2 \\
&= \frac{|\mathbf{x} - \mathbf{x}'|^2 v^2}{c^2} \cos^2 \theta' + \frac{|\mathbf{x} - \mathbf{x}'|^2 v^2}{c^2} \sin^2 \theta' \\
&= \frac{|\mathbf{x} - \mathbf{x}'|^2 v^2}{c^2} (\cos^2 \theta' + \sin^2 \theta') = \frac{|\mathbf{x} - \mathbf{x}'|^2 v^2}{c^2}
\end{aligned} \tag{8.95}$$

we find that

$$\left(\frac{(\mathbf{x} - \mathbf{x}') \cdot \mathbf{v}}{c}\right)^2 = \frac{|\mathbf{x} - \mathbf{x}'|^2 v^2}{c^2} - \left(\frac{(\mathbf{x} - \mathbf{x}') \times \mathbf{v}}{c}\right)^2 \quad (8.96)$$

Furthermore, from equation (8.87) on page 128, we obtain the following identity:

$$(\mathbf{x} - \mathbf{x}') \times \mathbf{v} = (\mathbf{x} - \mathbf{x}_0) \times \mathbf{v} \quad (8.97)$$

which, when inserted into equation (8.96) above, yields the relation

$$\left(\frac{(\mathbf{x} - \mathbf{x}') \cdot \mathbf{v}}{c}\right)^2 = \frac{|\mathbf{x} - \mathbf{x}'|^2 v^2}{c^2} - \left(\frac{(\mathbf{x} - \mathbf{x}_0) \times \mathbf{v}}{c}\right)^2 \quad (8.98)$$

Inserting the above into expression (8.94) on the preceding page for s^2 , this expression becomes

$$\begin{aligned} s^2 &= |\mathbf{x} - \mathbf{x}'|^2 - 2|\mathbf{x} - \mathbf{x}'| \frac{(\mathbf{x} - \mathbf{x}') \cdot \mathbf{v}}{c} + \frac{|\mathbf{x} - \mathbf{x}'|^2 v^2}{c^2} - \left(\frac{(\mathbf{x} - \mathbf{x}_0) \times \mathbf{v}}{c}\right)^2 \\ &= \left(\mathbf{x} - \mathbf{x}' - \frac{|\mathbf{x} - \mathbf{x}'| \mathbf{v}}{c}\right)^2 - \left(\frac{(\mathbf{x} - \mathbf{x}_0) \times \mathbf{v}}{c}\right)^2 \\ &= (\mathbf{x} - \mathbf{x}_0(t))^2 - \left(\frac{(\mathbf{x} - \mathbf{x}_0(t)) \times \mathbf{v}}{c}\right)^2 \\ &\equiv |\mathbf{x} - \mathbf{x}_0|^2 - \left(\frac{(\mathbf{x} - \mathbf{x}_0) \times \mathbf{v}}{c}\right)^2 \end{aligned} \quad (8.99)$$

where in the penultimate step we used equation (8.87) on page 128.

What we have just demonstrated is that, in the case the particle velocity at time t can be calculated or projected, the retarded distance s in the Liénard-Wiechert potentials (8.68) can be expressed in terms of the virtual simultaneous coordinate $\mathbf{x}_0(t)$, *viz.*, the point at which the particle will have arrived at time t , *i.e.*, when we obtain the first knowledge of its existence at the source point \mathbf{x}' at the retarded time t' , and in the field coordinate $\mathbf{x}(t)$, where we make our observations. We have, in other words, shown that all quantities in the definition of s , and hence s itself, can, when the motion of the charge is somehow known, be expressed in terms of the time t alone. *I.e.*, in this special case we are able to express the retarded relative distance as $s = s(t, \mathbf{x})$ and we do not have to involve the retarded time t' or transformed differential operators in our calculations.

Taking the square root of both sides of equation (8.99), we obtain the following alternative final expressions for the retarded relative distance s in terms of the

charge's virtual simultaneous coordinate $\mathbf{x}_0(t)$:

$$s(t, \mathbf{x}) = \sqrt{|\mathbf{x} - \mathbf{x}_0|^2 - \left(\frac{(\mathbf{x} - \mathbf{x}_0) \times \mathbf{v}}{c}\right)^2} \quad (8.100a)$$

$$= |\mathbf{x} - \mathbf{x}_0| \sqrt{1 - \frac{v^2}{c^2} \sin^2 \theta_0} \quad (8.100b)$$

$$= \sqrt{|\mathbf{x} - \mathbf{x}_0|^2 \left(1 - \frac{v^2}{c^2}\right) - \left(\frac{(\mathbf{x} - \mathbf{x}_0) \cdot \mathbf{v}}{c}\right)^2} \quad (8.100c)$$

Using equation (8.100c) above and standard vector analytic formulae, we obtain

$$\begin{aligned} \nabla s^2 &= \nabla \left[|\mathbf{x} - \mathbf{x}_0|^2 \left(1 - \frac{v^2}{c^2}\right) - \left(\frac{(\mathbf{x} - \mathbf{x}_0) \cdot \mathbf{v}}{c}\right)^2 \right] \\ &= 2 \left[(\mathbf{x} - \mathbf{x}_0) \left(1 - \frac{v^2}{c^2}\right) + \frac{\mathbf{v}\mathbf{v}}{c^2} \cdot (\mathbf{x} - \mathbf{x}_0) \right] \\ &= 2 \left[(\mathbf{x} - \mathbf{x}_0) + \frac{\mathbf{v}}{c} \times \left(\frac{\mathbf{v}}{c} \times (\mathbf{x} - \mathbf{x}_0)\right) \right] \end{aligned} \quad (8.101)$$

which we shall use in the following example of a uniformly, unaccelerated motion of the charge.

▷ THE FIELDS FROM A UNIFORMLY MOVING CHARGE

EXAMPLE 8.1

In the special case of uniform motion, the localised charge moves in a field-free, isolated space and we know that it will not be affected by any external forces. It will therefore move uniformly in a straight line with the constant velocity \mathbf{v} . This gives us the possibility to extrapolate its position at the observation time, $\mathbf{x}'(t)$, from its position at the retarded time, $\mathbf{x}'(t')$. Since the particle is not accelerated, $\dot{\mathbf{v}} \equiv \mathbf{0}$, the virtual simultaneous coordinate \mathbf{x}_0 will be identical to the actual *simultaneous coordinate* of the particle at time t , i.e., $\mathbf{x}_0(t) = \mathbf{x}'(t)$. As depicted in figure 8.3 on page 124, the angle between $\mathbf{x} - \mathbf{x}_0$ and \mathbf{v} is θ_0 while then angle between $\mathbf{x} - \mathbf{x}'$ and \mathbf{v} is θ' .

We note that in the case of uniform velocity \mathbf{v} , time and space derivatives are closely related in the following way when they operate on functions of $\mathbf{x}(t)$:

$$\frac{\partial}{\partial t} \rightarrow -\mathbf{v} \cdot \nabla \quad (8.102)$$

Hence, the \mathbf{E} and \mathbf{B} fields can be obtained from formulae (8.71) on page 124, with the

potentials given by equations (8.68) on page 123 as follows:

$$\begin{aligned}
 \mathbf{E} &= -\nabla\phi - \frac{\partial\mathbf{A}}{\partial t} = -\nabla\phi - \frac{1}{c^2} \frac{\partial\mathbf{v}\phi}{\partial t} = -\nabla\phi - \frac{\mathbf{v}}{c^2} \frac{\partial\phi}{\partial t} \\
 &= -\nabla\phi + \frac{\mathbf{v}}{c} \left(\frac{\mathbf{v}}{c} \cdot \nabla\phi \right) = -\left(1 - \frac{\mathbf{v}\mathbf{v}}{c^2}\right) \nabla\phi \\
 &= \left(\frac{\mathbf{v}\mathbf{v}}{c^2} - \mathbf{1}\right) \cdot \nabla\phi
 \end{aligned} \tag{8.103a}$$

$$\begin{aligned}
 \mathbf{B} &= \nabla \times \mathbf{A} = \nabla \times \left(\frac{\mathbf{v}}{c^2} \phi \right) = \nabla\phi \times \frac{\mathbf{v}}{c^2} = -\frac{\mathbf{v}}{c^2} \times \nabla\phi \\
 &= \frac{\mathbf{v}}{c^2} \times \left[\left(\frac{\mathbf{v}}{c} \cdot \nabla\phi \right) \frac{\mathbf{v}}{c} - \nabla\phi \right] = \frac{\mathbf{v}}{c^2} \times \left(\frac{\mathbf{v}\mathbf{v}}{c^2} - \mathbf{1} \right) \cdot \nabla\phi \\
 &= \frac{\mathbf{v}}{c^2} \times \mathbf{E}
 \end{aligned} \tag{8.103b}$$

Here $\mathbf{1} = \hat{x}_i \hat{x}_i$ is the unit dyad and we used the fact that $\mathbf{v} \times \mathbf{v} \equiv 0$. What remains is just to express $\nabla\phi$ in quantities evaluated at t and \mathbf{x} .

From equation (8.68a) on page 123 and equation (8.101) on the preceding page we find that

$$\begin{aligned}
 \nabla\phi &= \frac{q'}{4\pi\epsilon_0} \nabla \left(\frac{1}{s} \right) = -\frac{q'}{8\pi\epsilon_0 s^3} \nabla s^2 \\
 &= -\frac{q'}{4\pi\epsilon_0 s^3} \left[(\mathbf{x} - \mathbf{x}_0) + \frac{\mathbf{v}}{c} \times \left(\frac{\mathbf{v}}{c} \times (\mathbf{x} - \mathbf{x}_0) \right) \right]
 \end{aligned} \tag{8.104}$$

When this expression for $\nabla\phi$ is inserted into equation (8.103a) above, the following result

$$\begin{aligned}
 \mathbf{E}(t, \mathbf{x}) &= \left(\frac{\mathbf{v}\mathbf{v}}{c^2} - \mathbf{1} \right) \cdot \nabla\phi = -\frac{q'}{8\pi\epsilon_0 s^3} \left(\frac{\mathbf{v}\mathbf{v}}{c^2} - \mathbf{1} \right) \cdot \nabla s^2 \\
 &= \frac{q'}{4\pi\epsilon_0 s^3} \left\{ (\mathbf{x} - \mathbf{x}_0) + \frac{\mathbf{v}}{c} \times \left(\frac{\mathbf{v}}{c} \times (\mathbf{x} - \mathbf{x}_0) \right) \right. \\
 &\quad \left. - \frac{\mathbf{v}}{c} \left(\frac{\mathbf{v}}{c} \cdot (\mathbf{x} - \mathbf{x}_0) \right) - \frac{\mathbf{v}\mathbf{v}}{c^2} \cdot \left[\frac{\mathbf{v}}{c} \times \left(\frac{\mathbf{v}}{c} \times (\mathbf{x} - \mathbf{x}_0) \right) \right] \right\} \\
 &= \frac{q'}{4\pi\epsilon_0 s^3} \left[(\mathbf{x} - \mathbf{x}_0) + \frac{\mathbf{v}}{c} \left(\frac{\mathbf{v}}{c} \cdot (\mathbf{x} - \mathbf{x}_0) \right) - (\mathbf{x} - \mathbf{x}_0) \frac{v^2}{c^2} \right. \\
 &\quad \left. - \frac{\mathbf{v}}{c} \left(\frac{\mathbf{v}}{c} \cdot (\mathbf{x} - \mathbf{x}_0) \right) \right] \\
 &= \frac{q'}{4\pi\epsilon_0 s^3} (\mathbf{x} - \mathbf{x}_0) \left(1 - \frac{v^2}{c^2} \right)
 \end{aligned} \tag{8.105}$$

follows. Of course, the same result also follows from equation (8.88) on page 128 with $\dot{\mathbf{v}} \equiv \mathbf{0}$ inserted.

From equation (8.105) above we conclude that \mathbf{E} is directed along the vector from the simultaneous coordinate $\mathbf{x}_0(t)$ to the field (observation) coordinate $\mathbf{x}(t)$. In a similar way, the magnetic field can be calculated and one finds that

$$\mathbf{B}(t, \mathbf{x}) = \frac{\mu_0 q'}{4\pi s^3} \left(1 - \frac{v^2}{c^2}\right) \mathbf{v} \times (\mathbf{x} - \mathbf{x}_0) = \frac{1}{c^2} \mathbf{v} \times \mathbf{E} \quad (8.106)$$

From these explicit formulae for the \mathbf{E} and \mathbf{B} fields we can discern the following cases:

1. $v \rightarrow 0 \Rightarrow \mathbf{E}$ goes over into the Coulomb field $\mathbf{E}^{\text{Coulomb}}$
2. $v \rightarrow 0 \Rightarrow \mathbf{B}$ goes over into the Biot-Savart field
3. $v \rightarrow c \Rightarrow \mathbf{E}$ becomes dependent on θ_0
4. $v \rightarrow c, \sin \theta_0 \approx 0 \Rightarrow \mathbf{E} \rightarrow (1 - v^2/c^2) \mathbf{E}^{\text{Coulomb}}$
5. $v \rightarrow c, \sin \theta_0 \approx 1 \Rightarrow \mathbf{E} \rightarrow (1 - v^2/c^2)^{-1/2} \mathbf{E}^{\text{Coulomb}}$

END OF EXAMPLE 8.1 ◀

▷ THE CONVECTION POTENTIAL AND THE CONVECTION FORCE

EXAMPLE 8.2

Let us consider in more detail the treatment of the radiation from a uniformly moving rigid charge distribution.

If we return to the original definition of the potentials and the inhomogeneous wave equation, formula (3.19) on page 38, for a generic potential component $\Psi(t, \mathbf{x})$ and a generic source component $f(t, \mathbf{x})$,

$$\square^2 \Psi(t, \mathbf{x}) = \left(\frac{1}{c^2} \frac{\partial^2}{\partial t^2} - \nabla^2 \right) \Psi(t, \mathbf{x}) = f(t, \mathbf{x}) \quad (8.107)$$

we find that under the assumption that $\mathbf{v} = v \hat{\mathbf{x}}_1$, this equation can be written

$$\left(1 - \frac{v^2}{c^2}\right) \frac{\partial^2 \Psi}{\partial x_1^2} + \frac{\partial^2 \Psi}{\partial x_2^2} + \frac{\partial^2 \Psi}{\partial x_3^2} = -f(\mathbf{x}) \quad (8.108)$$

i.e., in a time-independent form. Transforming

$$\xi_1 = \frac{x_1}{\sqrt{1 - v^2/c^2}} \quad (8.109a)$$

$$\xi_2 = x_2 \quad (8.109b)$$

$$\xi_3 = x_3 \quad (8.109c)$$

and introducing the vectorial nabla operator in ξ space, $\nabla_\xi \stackrel{\text{def}}{=} (\partial/\partial \xi_1, \partial/\partial \xi_2, \partial/\partial \xi_3)$,

the time-independent equation (8.108) reduces to an ordinary *Poisson equation*

$$\nabla_{\xi}^2 \Psi(\xi) = -f(\sqrt{1 - v^2/c^2} \xi_1, \xi_2, \xi_3) \equiv -f(\xi) \quad (8.110)$$

in this space. This equation has the well-known Coulomb potential solution

$$\Psi(\xi) = \frac{1}{4\pi} \int_V \frac{f(\xi')}{|\xi - \xi'|} d^3\xi' \quad (8.111)$$

After inverse transformation back to the original coordinates, this becomes

$$\Psi(\mathbf{x}) = \frac{1}{4\pi} \int_V \frac{f(\mathbf{x}')}{s} d^3x' \quad (8.112)$$

where, in the denominator,

$$s = \left[(x_1 - x'_1)^2 + \left(1 - \frac{v^2}{c^2}\right) [(x_2 - x'_2)^2 + (x_3 - x'_3)^2] \right]^{\frac{1}{2}} \quad (8.113)$$

Applying this to the explicit scalar and vector potential components, realising that for a rigid charge distribution ρ moving with velocity \mathbf{v} the current is given by $\mathbf{j} = \rho\mathbf{v}$, we obtain

$$\phi(t, \mathbf{x}) = \frac{1}{4\pi\epsilon_0} \int_V \frac{\rho(\mathbf{x}')}{s} d^3x' \quad (8.114a)$$

$$\mathbf{A}(t, \mathbf{x}) = \frac{1}{4\pi\epsilon_0 c^2} \int_V \frac{\mathbf{v}\rho(\mathbf{x}')}{s} d^3x' = \frac{\mathbf{v}}{c^2} \phi(t, \mathbf{x}) \quad (8.114b)$$

For a localised charge where $\int \rho d^3x' = q'$, these expressions reduce to

$$\phi(t, \mathbf{x}) = \frac{q'}{4\pi\epsilon_0 s} \quad (8.115a)$$

$$\mathbf{A}(t, \mathbf{x}) = \frac{q'\mathbf{v}}{4\pi\epsilon_0 c^2 s} \quad (8.115b)$$

which we recognise as the *Liénard-Wiechert potentials*; cf. equations (8.68) on page 123. We notice, however, that the derivation here, based on a mathematical technique which in fact is a *Lorentz transformation*, is of more general validity than the one leading to equations (8.68) on page 123.

Let us now consider the action of the fields produced from a moving, rigid charge distribution represented by q' moving with velocity \mathbf{v} , on a charged particle q , also moving with velocity \mathbf{v} . This force is given by the *Lorentz force*

$$\mathbf{F} = q(\mathbf{E} + \mathbf{v} \times \mathbf{B}) \quad (8.116)$$

This means that we can rewrite expression (8.116) above as

$$\mathbf{F} = q \left[\mathbf{E} + \mathbf{v} \times \left(\frac{\mathbf{v}}{c^2} \times \mathbf{E} \right) \right] = q \left[\left(\frac{\mathbf{v}}{c} \cdot \nabla \phi \right) \frac{\mathbf{v}}{c} - \nabla \phi - \frac{\mathbf{v}}{c} \times \left(\frac{\mathbf{v}}{c} \times \nabla \phi \right) \right] \quad (8.117)$$

Applying the “bac-cab” rule, formula (F.56) on page 167, on the last term yields

$$\frac{\mathbf{v}}{c} \times \left(\frac{\mathbf{v}}{c} \times \nabla \phi \right) = \left(\frac{\mathbf{v}}{c} \cdot \nabla \phi \right) \frac{\mathbf{v}}{c} - \frac{v^2}{c^2} \nabla \phi \quad (8.118)$$

which means that we can write

$$\mathbf{F} = -q \nabla \psi \quad (8.119)$$

where

$$\psi = \left(1 - \frac{v^2}{c^2} \right) \phi \quad (8.120)$$

The scalar function ψ is called the *convection potential* or the *Heaviside potential*. When the rigid charge distribution is well localised so that we can use the potentials (8.115) the convection potential becomes

$$\psi = \left(1 - \frac{v^2}{c^2} \right) \frac{q'}{4\pi\epsilon_0 s} \quad (8.121)$$

The convection potential from a point charge is constant on flattened ellipsoids of revolution, defined through equation (8.113) on the preceding page as

$$\begin{aligned} & \left(\frac{x_1 - x'_1}{\sqrt{1 - v^2/c^2}} \right)^2 + (x_2 - x'_2)^2 + (x_3 - x'_3)^2 \\ & = \gamma^2 (x_1 - x'_1)^2 + (x_2 - x'_2)^2 + (x_3 - x'_3)^2 = \text{Const} \end{aligned} \quad (8.122)$$

These Heaviside ellipsoids are equipotential surfaces, and since the force is proportional to the gradient of ψ , which means that it is perpendicular to the ellipsoid surface, the force between the two charges is in general *not* directed along the line which connects the charges. A consequence of this is that a system consisting of two co-moving charges connected with a rigid bar, will experience a torque. This is the idea behind the Trouton-Noble experiment, aimed at measuring the *absolute speed* of the earth or the galaxy. The negative outcome of this experiment is explained by the special theory of relativity which postulates that mechanical laws follow the same rules as electromagnetic laws, so that a compensating torque appears due to mechanical stresses within the charge-bar system.

END OF EXAMPLE 8.2◁

Radiation for small velocities

If the charge moves at such low speeds that $v/c \ll 1$, formula (8.69) on page 123 simplifies to

$$s = |\mathbf{x} - \mathbf{x}'| - \frac{(\mathbf{x} - \mathbf{x}') \cdot \mathbf{v}}{c} \approx |\mathbf{x} - \mathbf{x}'|, \quad v \ll c \quad (8.123)$$

and formula (8.87) on page 128

$$\mathbf{x} - \mathbf{x}_0 = (\mathbf{x} - \mathbf{x}') - \frac{|\mathbf{x} - \mathbf{x}'| \mathbf{v}}{c} \approx \mathbf{x} - \mathbf{x}', \quad v \ll c \quad (8.124)$$

so that the radiation field equation (8.92) on page 129 can be approximated by

$$\mathbf{E}^{\text{rad}}(t, \mathbf{x}) = \frac{q'}{4\pi\epsilon_0 c^2 |\mathbf{x} - \mathbf{x}'|^3} (\mathbf{x} - \mathbf{x}') \times [(\mathbf{x} - \mathbf{x}') \times \dot{\mathbf{v}}], \quad v \ll c \quad (8.125)$$

from which we obtain, with the use of formula (8.91) on page 129, the magnetic field

$$\mathbf{B}^{\text{rad}}(t, \mathbf{x}) = \frac{q'}{4\pi c^3 |\mathbf{x} - \mathbf{x}'|^2} [\dot{\mathbf{v}} \times (\mathbf{x} - \mathbf{x}')], \quad v \ll c \quad (8.126)$$

It is interesting to note the close correspondence which exists between the non-relativistic fields (8.125) and (8.126) and the electric dipole field equations (8.47) on page 118 if we introduce

$$\mathbf{p} = q' \mathbf{x}'(t') \quad (8.127)$$

and at the same time make the transitions

$$q' \dot{\mathbf{v}} = \ddot{\mathbf{p}} \rightarrow -\omega^2 \mathbf{p}_\omega \quad (8.128a)$$

$$\mathbf{x} - \mathbf{x}' = \mathbf{x} - \mathbf{x}_0 \quad (8.128b)$$

The power flux in the far zone is described by the Poynting vector as a function of \mathbf{E}^{rad} and \mathbf{B}^{rad} . We use the close correspondence with the dipole case to find that it becomes

$$\mathbf{S} = \frac{\mu_0 q'^2 (\dot{\mathbf{v}})^2}{16\pi^2 c |\mathbf{x} - \mathbf{x}'|^2} \sin^2 \theta \frac{\mathbf{x} - \mathbf{x}'}{|\mathbf{x} - \mathbf{x}'|} \quad (8.129)$$

where θ is the angle between $\dot{\mathbf{v}}$ and $\mathbf{x} - \mathbf{x}_0$. The total radiated power (integrated over a closed spherical surface) becomes

$$P = \frac{\mu_0 q'^2 (\dot{\mathbf{v}})^2}{6\pi c} = \frac{q'^2 \dot{v}^2}{6\pi\epsilon_0 c^3} \quad (8.130)$$

which is the *Larmor formula for radiated power* from an accelerated charge. Note that here we are treating a charge with $v \ll c$ but otherwise *totally unspecified motion* while we compare with formulae derived for a *stationary oscillating dipole*. The electric and magnetic fields, equation (8.125) above and equation (8.126), respectively, and the expressions for the Poynting flux and power derived from them, are here *instantaneous* values, dependent on the instantaneous position of the charge at $\mathbf{x}'(t')$. The angular distribution is that which is “frozen” to the point from which the energy is radiated.

8.5.3 Bremsstrahlung

An important special case of radiation is when the velocity \mathbf{v} and the acceleration $\dot{\mathbf{v}}$ are collinear (parallel or anti-parallel) so that $\mathbf{v} \times \dot{\mathbf{v}} = \mathbf{0}$. This condition (for an arbitrary magnitude of \mathbf{v}) inserted into expression (8.92) on page 129 for the radiation field, yields

$$\mathbf{E}^{\text{rad}}(t, \mathbf{x}) = \frac{q'}{4\pi\epsilon_0 c^2 s^3} (\mathbf{x} - \mathbf{x}') \times [(\mathbf{x} - \mathbf{x}') \times \dot{\mathbf{v}}], \quad \mathbf{v} \parallel \dot{\mathbf{v}} \quad (8.131)$$

from which we obtain, with the use of formula (8.91) on page 129, the magnetic field

$$\mathbf{B}^{\text{rad}}(t, \mathbf{x}) = \frac{q' |\mathbf{x} - \mathbf{x}'|}{4\pi\epsilon_0 c^3 s^3} [\dot{\mathbf{v}} \times (\mathbf{x} - \mathbf{x}')], \quad \mathbf{v} \parallel \dot{\mathbf{v}} \quad (8.132)$$

The difference between this case and the previous case of $v \ll c$ is that the approximate expression (8.123) on page 135 for s is no longer valid; we must instead use the correct expression (8.69) on page 123. The angular distribution of the power flux (Poynting vector) therefore becomes

$$\mathbf{S} = \frac{\mu_0 q'^2 \dot{v}^2}{16\pi^2 c |\mathbf{x} - \mathbf{x}'|^2} \frac{\sin^2 \theta}{\left(1 - \frac{v}{c} \cos \theta\right)^6} \frac{\mathbf{x} - \mathbf{x}'}{|\mathbf{x} - \mathbf{x}'|} \quad (8.133)$$

It is interesting to note that the magnitudes of the electric and magnetic fields are the same whether \mathbf{v} and $\dot{\mathbf{v}}$ are parallel or anti-parallel.

We must be careful when we compute the energy (\mathbf{S} integrated over time). The Poynting vector is related to the time t when it is measured and to a *fixed* surface in space. The radiated power into a solid angle element $d\Omega$, measured relative to the particle's retarded position, is given by the formula

$$\frac{dU^{\text{rad}}(\theta)}{dt} d\Omega = \mathbf{S} \cdot (\mathbf{x} - \mathbf{x}') |\mathbf{x} - \mathbf{x}'| d\Omega = \frac{\mu_0 q'^2 \dot{v}^2}{16\pi^2 c} \frac{\sin^2 \theta}{\left(1 - \frac{v}{c} \cos \theta\right)^6} d\Omega \quad (8.134)$$

On the other hand, the radiation loss due to radiation from the charge at retarded time t' :

$$\frac{dU^{\text{rad}}}{dt'} d\Omega = \frac{dU^{\text{rad}}}{dt} \left(\frac{\partial t}{\partial t'} \right)_{\mathbf{x}} d\Omega \quad (8.135)$$

Using formula (8.79) on page 126, we obtain

$$\frac{dU^{\text{rad}}}{dt'} d\Omega = \frac{dU^{\text{rad}}}{dt} \frac{s}{|\mathbf{x} - \mathbf{x}'|} d\Omega = \mathbf{S} \cdot (\mathbf{x} - \mathbf{x}') s d\Omega \quad (8.136)$$

Inserting equation (8.133) for \mathbf{S} into (8.136), we obtain the explicit expression

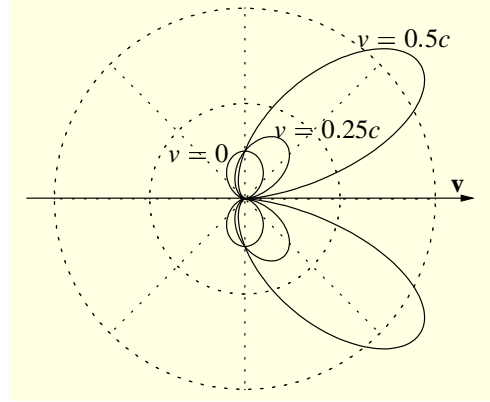


Figure 8.4. Polar diagram of the energy loss angular distribution factor $\sin^2 \theta / (1 - v \cos \theta / c)^5$ during bremsstrahlung for particle speeds $v = 0$, $v = 0.25c$, and $v = 0.5c$.

for the energy loss due to radiation evaluated at the retarded time

$$\frac{dU^{\text{rad}}(\theta)}{dt'} d\Omega = \frac{\mu_0 q^2 \dot{v}^2}{16\pi^2 c} \frac{\sin^2 \theta}{(1 - \frac{v}{c} \cos \theta)^5} d\Omega \quad (8.137)$$

The angular factors of this expression, for three different particle speeds, are plotted in Figure 8.4.

Comparing expression (8.134) on the previous page with expression (8.137), we see that they differ by a factor $1 - v \cos \theta / c$ which comes from the extra factor $s / |\mathbf{x} - \mathbf{x}'|$ introduced in (8.136). Let us explain this in geometrical terms.

During the interval $(t', t' + dt')$ and within the solid angle element $d\Omega$ the particle radiates an energy $[dU^{\text{rad}}(\theta) / dt'] dt' d\Omega$. As shown in Figure 8.5 this energy is at time t located between two spheres, one outer with its origin in $\mathbf{x}'_1(t')$ and one inner with its origin in $\mathbf{x}'_1(t' + dt') = \mathbf{x}'_1(t') + \mathbf{v} dt'$ and radius $c[t - (t' + dt')] = c(t - t' - dt')$.

From Figure 8.5 we see that the volume element subtending the solid angle element

$$d\Omega = \frac{dS}{|\mathbf{x} - \mathbf{x}'_2|^2} \quad (8.138)$$

is

$$d^3x = dS dr = |\mathbf{x} - \mathbf{x}'_2|^2 d\Omega dr \quad (8.139)$$

Here, dr denotes the differential distance between the two spheres and can be eval-

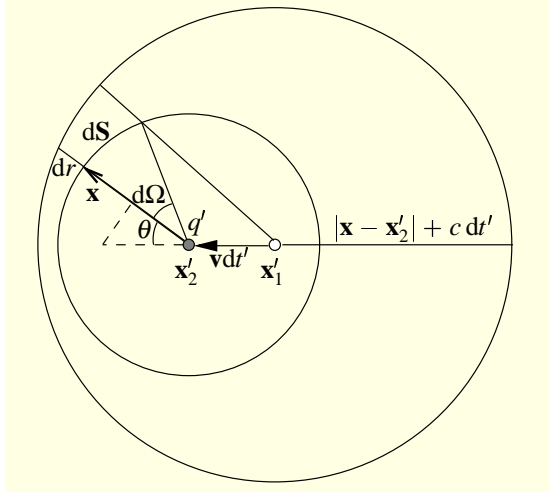


Figure 8.5. Location of radiation between two spheres as the charge moves with velocity \mathbf{v} from \mathbf{x}'_1 to \mathbf{x}'_2 during the time interval $(t', t' + dt')$. The observation point (field point) is at the fixed location \mathbf{x} .

uated in the following way

$$\begin{aligned} dr &= |\mathbf{x} - \mathbf{x}'_2| + c dt' - \underbrace{\frac{\mathbf{x} - \mathbf{x}'_2}{|\mathbf{x} - \mathbf{x}'_2|} \cdot \mathbf{v}}_{v \cos \theta} dt' - |\mathbf{x} - \mathbf{x}'_2| \\ &= \left(c - \frac{\mathbf{x} - \mathbf{x}'_2}{|\mathbf{x} - \mathbf{x}'_2|} \cdot \mathbf{v} \right) dt' = \frac{cs}{|\mathbf{x} - \mathbf{x}'_2|} dt' \end{aligned} \quad (8.140)$$

where formula (8.69) on page 123 was used in the last step. Hence, the volume element under consideration is

$$d^3x = dS dr = \frac{s}{|\mathbf{x} - \mathbf{x}'_2|} dS c dt' \quad (8.141)$$

We see that the energy which is radiated per unit solid angle during the time interval $(t', t' + dt')$ is located in a volume element whose size is θ dependent. This explains the difference between expression (8.134) on page 137 and expression (8.137) on the preceding page.

Let the radiated energy, integrated over Ω , be denoted \tilde{U}^{rad} . After tedious, but relatively straightforward integration of formula (8.137) on the facing page, one

obtains

$$\frac{d\tilde{U}^{\text{rad}}}{dt'} = \frac{\mu_0 q'^2 \dot{v}^2}{6\pi c} \frac{1}{\left(1 - \frac{v^2}{c^2}\right)^3} \quad (8.142)$$

If we know $\mathbf{v}(t')$, we can integrate this expression over t' and obtain the total energy radiated during the acceleration or deceleration of the particle. This way we obtain a classical picture of *bremstrahlung* (*braking radiation*). Often, an atomistic treatment is required for an acceptable result.

EXAMPLE 8.3

▷ BREMSSTRAHLUNG FOR LOW SPEEDS AND SHORT ACCELERATION TIMES

Calculate the bremsstrahlung when a charged particle, moving at a non-relativistic speed, is accelerated or decelerated during an infinitely short time interval.

We approximate the velocity change at time $t' = t_0$ by a delta function:

$$\dot{\mathbf{v}}(t') = \Delta\mathbf{v} \delta(t' - t_0) \quad (8.143)$$

which means that

$$\Delta\mathbf{v} = \int_{-\infty}^{\infty} \dot{\mathbf{v}} dt \quad (8.144)$$

Also, we assume $v/c \ll 1$ so that, according to formula (8.69) on page 123,

$$s \approx |\mathbf{x} - \mathbf{x}'| \quad (8.145)$$

and, according to formula (8.87) on page 128,

$$\mathbf{x} - \mathbf{x}_0 \approx \mathbf{x} - \mathbf{x}' \quad (8.146)$$

From the general expression (8.91) on page 129 we conclude that $\mathbf{E} \perp \mathbf{B}$ and that it suffices to consider $E \equiv |\mathbf{E}^{\text{rad}}|$. According to the “bremsstrahlung expression” for \mathbf{E}^{rad} , equation (8.131) on page 137,

$$E = \frac{q' \sin \theta}{4\pi\epsilon_0 c^2 |\mathbf{x} - \mathbf{x}'|} \Delta v \delta(t' - t_0) \quad (8.147)$$

In this simple case $B \equiv |\mathbf{B}^{\text{rad}}|$ is given by

$$B = \frac{E}{c} \quad (8.148)$$

Fourier transforming expression (8.147) above for E is trivial, yielding

$$E_\omega = \frac{q' \sin \theta}{8\pi^2 \epsilon_0 c^2 |\mathbf{x} - \mathbf{x}'|} \Delta v e^{i\omega t_0} \quad (8.149)$$

We note that the magnitude of this Fourier component is independent of ω . This is a consequence of the infinitely short “impulsive step” $\delta(t' - t_0)$ in the time domain which produces an infinite spectrum in the frequency domain.

The total radiation energy is given by the expression

$$\begin{aligned}
\tilde{U}^{\text{rad}} &= \int \frac{d\tilde{U}^{\text{rad}}}{dt'} dt' = \int_{-\infty}^{\infty} \int_S \left(\mathbf{E} \times \frac{\mathbf{B}}{\mu_0} \right) \cdot d\mathbf{S} dt' \\
&= \frac{1}{\mu_0} \int_S \int_{-\infty}^{\infty} EB dt' d^2x = \frac{1}{\mu_0 c} \int_S \int_{-\infty}^{\infty} E^2 dt' d^2x \\
&= \varepsilon_0 c \int_S \int_{-\infty}^{\infty} E^2 dt' d^2x
\end{aligned} \tag{8.150}$$

According to *Parseval's identity* [cf. equation (8.15) on page 111] the following equality holds:

$$\int_{-\infty}^{\infty} E^2 dt' = 4\pi \int_0^{\infty} |E_\omega|^2 d\omega \tag{8.151}$$

which means that the radiated energy in the frequency interval $(\omega, \omega + d\omega)$ is

$$\tilde{U}_\omega^{\text{rad}} d\omega = 4\pi\varepsilon_0 c \left(\int_S |E_\omega|^2 d^2x \right) d\omega \tag{8.152}$$

For our infinite spectrum, equation (8.149) on the preceding page, we obtain

$$\begin{aligned}
\tilde{U}_\omega^{\text{rad}} d\omega &= \frac{q'^2 (\Delta v)^2}{16\pi^3 \varepsilon_0 c^3} \int_S \frac{\sin^2 \theta}{|\mathbf{x} - \mathbf{x}'|^2} d^2x d\omega \\
&= \frac{q'^2 (\Delta v)^2}{16\pi^3 \varepsilon_0 c^3} \int_0^{2\pi} d\varphi \int_0^\pi \sin^2 \theta \sin \theta d\theta d\omega \\
&= \frac{q'^2}{3\pi\varepsilon_0 c} \left(\frac{\Delta v}{c} \right)^2 \frac{d\omega}{2\pi}
\end{aligned} \tag{8.153}$$

We see that the energy spectrum $\tilde{U}_\omega^{\text{rad}}$ is independent of frequency ω . This means that if we integrate it over all frequencies $\omega \in [0, \infty]$, a divergent integral would result.

In reality, all spectra have finite widths, with an upper *cutoff* limit set by the quantum condition

$$\hbar\omega = \frac{1}{2}m(\Delta v)^2 \tag{8.154}$$

which expresses that the highest possible frequency in the spectrum is that for which all kinetic energy difference has gone into one single *field quantum (photon)* with energy $\hbar\omega$. If we adopt the picture that the total energy is quantised in terms of N_ω photons radiated during the process, we find that

$$\frac{\tilde{U}_\omega^{\text{rad}} d\omega}{\hbar\omega} = dN_\omega \tag{8.155}$$

or, for an electron where $q' = -|e|$, where e is the elementary charge,

$$dN_\omega = \frac{e^2}{4\pi\varepsilon_0 \hbar c} \frac{2}{3\pi} \left(\frac{\Delta v}{c} \right)^2 \frac{d\omega}{\omega} \approx \frac{1}{137} \frac{2}{3\pi} \left(\frac{\Delta v}{c} \right)^2 \frac{d\omega}{\omega} \tag{8.156}$$

where we used the value of the *fine structure constant* $e^2/(4\pi\varepsilon_0 \hbar c) \approx 1/137$.

Even if the number of photons becomes infinite when $\omega \rightarrow 0$, these photons have negligi-

ble energies so that the total radiated energy is still finite.

END OF EXAMPLE 8.3◁

8.5.4 Cyclotron and synchrotron radiation

Formula (8.91) and formula (8.92) on page 129 for the magnetic field and the radiation part of the electric field are general, valid for any kind of motion of the localised charge. A very important special case is circular motion, *i.e.*, the case $\mathbf{v} \perp \dot{\mathbf{v}}$.

With the charged particle orbiting in the x_1x_2 plane as in Figure 8.6, an orbit radius a , and an angular frequency ω_0 , we obtain

$$\varphi(t') = \omega_0 t' \quad (8.157a)$$

$$\mathbf{x}'(t') = a[\hat{\mathbf{x}}_1 \cos \varphi(t') + \hat{\mathbf{x}}_2 \sin \varphi(t')] \quad (8.157b)$$

$$\mathbf{v}(t') = \dot{\mathbf{x}}'(t') = a\omega_0[-\hat{\mathbf{x}}_1 \sin \varphi(t') + \hat{\mathbf{x}}_2 \cos \varphi(t')] \quad (8.157c)$$

$$v = |\mathbf{v}| = a\omega_0 \quad (8.157d)$$

$$\dot{\mathbf{v}}(t') = \ddot{\mathbf{x}}'(t') = -a\omega_0^2[\hat{\mathbf{x}}_1 \cos \varphi(t') + \hat{\mathbf{x}}_2 \sin \varphi(t')] \quad (8.157e)$$

$$\dot{v} = |\dot{\mathbf{v}}| = a\omega_0^2 \quad (8.157f)$$

Because of the rotational symmetry we can, without loss of generality, rotate our coordinate system in around the x_3 axis so the relative vector $\mathbf{x} - \mathbf{x}'$ from the source point to an arbitrary field point always lies in the x_2x_3 plane, *i.e.*,

$$\mathbf{x} - \mathbf{x}' = |\mathbf{x} - \mathbf{x}'| (\hat{\mathbf{x}}_2 \sin \alpha + \hat{\mathbf{x}}_3 \cos \alpha) \quad (8.158)$$

where α is the angle between $\mathbf{x} - \mathbf{x}'$ and the normal to the plane of the particle orbit (see Figure 8.6). From the above expressions we obtain

$$(\mathbf{x} - \mathbf{x}') \cdot \mathbf{v} = |\mathbf{x} - \mathbf{x}'| v \sin \alpha \cos \varphi \quad (8.159a)$$

$$(\mathbf{x} - \mathbf{x}') \cdot \dot{\mathbf{v}} = -|\mathbf{x} - \mathbf{x}'| \dot{v} \sin \alpha \sin \varphi = |\mathbf{x} - \mathbf{x}'| \dot{v} \cos \theta \quad (8.159b)$$

where in the last step we simply used the definition of a scalar product and the fact that the angle between $\dot{\mathbf{v}}$ and $\mathbf{x} - \mathbf{x}'$ is θ .

The power flux is given by the Poynting vector, which, with the help of for-

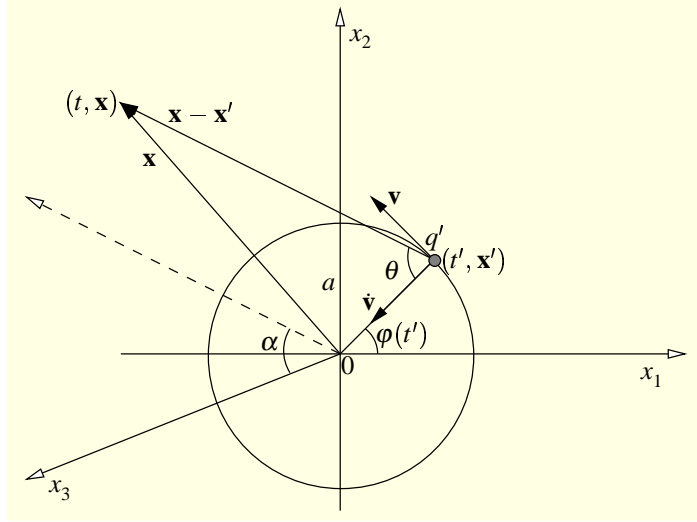


Figure 8.6. Coordinate system for the radiation from a charged particle at $\mathbf{x}'(t')$ in circular motion with velocity $\mathbf{v}(t')$ along the tangent and constant acceleration $\dot{\mathbf{v}}(t')$ toward the origin. The x_1x_2 axes are chosen so that the relative field point vector $\mathbf{x} - \mathbf{x}'$ makes an angle α with the x_3 axis which is normal to the plane of the orbital motion. The radius of the orbit is a .

mula (8.91) on page 129, can be written

$$\mathbf{S} = \frac{1}{\mu_0} (\mathbf{E} \times \mathbf{B}) = \frac{1}{c\mu_0} |\mathbf{E}|^2 \frac{\mathbf{x} - \mathbf{x}'}{|\mathbf{x} - \mathbf{x}'|} \quad (8.160)$$

Inserting this into equation (8.136) on page 137, we obtain

$$\frac{dU^{\text{rad}}(\alpha, \varphi)}{dt'} = \frac{|\mathbf{x} - \mathbf{x}'| s}{c\mu_0} |\mathbf{E}|^2 \quad (8.161)$$

where the retarded distance s is given by expression (8.69) on page 123. With the radiation part of the electric field, expression (8.92) on page 129, inserted, and using (8.159a) and (8.159b) on the preceding page, one finds, after some algebra, that

$$\frac{dU^{\text{rad}}(\alpha, \varphi)}{dt'} = \frac{\mu_0 q'^2 \dot{v}^2}{16\pi^2 c} \frac{\left(1 - \frac{v}{c} \sin \alpha \cos \varphi\right)^2 - \left(1 - \frac{v^2}{c^2}\right) \sin^2 \alpha \sin^2 \varphi}{\left(1 - \frac{v}{c} \sin \alpha \cos \varphi\right)^5} \quad (8.162)$$

The angles θ and φ vary in time during the rotation, so that θ refers to a *moving*

coordinate system. But we can parametrise the solid angle $d\Omega$ in the angle φ and the (fixed) angle α so that $d\Omega = \sin \alpha d\alpha d\varphi$. Integration of equation (8.162) on the preceding page over this $d\Omega$ gives, after some cumbersome algebra, the angular integrated expression

$$\frac{d\tilde{U}^{\text{rad}}}{dt'} = \frac{\mu_0 q'^2 \dot{v}^2}{6\pi c} \frac{1}{\left(1 - \frac{v^2}{c^2}\right)^2} \quad (8.163)$$

In equation (8.162) on the previous page, two limits are particularly interesting:

1. $v/c \ll 1$ which corresponds to *cyclotron radiation*.
2. $v/c \lesssim 1$ which corresponds to *synchrotron radiation*.

Cyclotron radiation

For a non-relativistic speed $v \ll c$, equation (8.162) on the preceding page reduces to

$$\frac{dU^{\text{rad}}(\alpha, \varphi)}{dt'} = \frac{\mu_0 q'^2 \dot{v}^2}{16\pi^2 c} (1 - \sin^2 \alpha \sin^2 \varphi) \quad (8.164)$$

But, according to equation (8.159b) on page 142

$$\sin^2 \alpha \sin^2 \varphi = \cos^2 \theta \quad (8.165)$$

where θ is defined in Figure 8.6. This means that we can write

$$\frac{dU^{\text{rad}}(\theta)}{dt'} = \frac{\mu_0 q'^2 \dot{v}^2}{16\pi^2 c} (1 - \cos^2 \theta) = \frac{\mu_0 q'^2 \dot{v}^2}{16\pi^2 c} \sin^2 \theta \quad (8.166)$$

Consequently, a fixed observer near the orbit plane will observe cyclotron radiation twice per revolution in the form of two equally broad pulses of radiation with alternating polarisation.

Synchrotron radiation

When the particle is relativistic, $v \lesssim c$, the denominator in equation (8.162) on the preceding page becomes very small if $\sin \alpha \cos \varphi \approx 1$, which defines the forward direction of the particle motion ($\alpha \approx \pi/2$, $\varphi \approx 0$). equation (8.162) on the

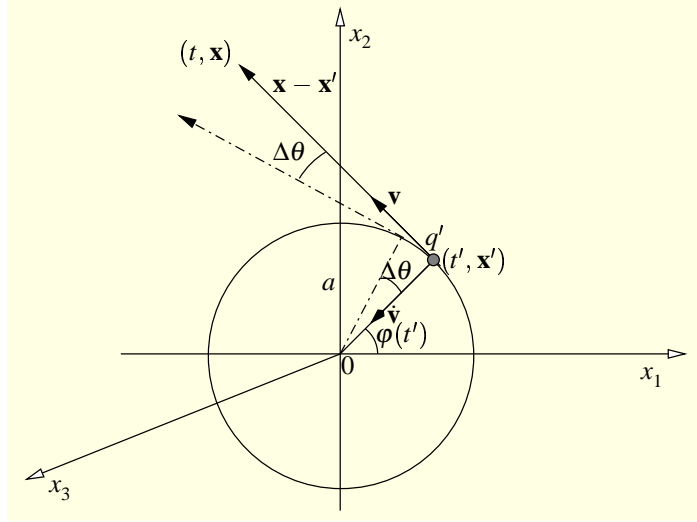


Figure 8.7. When the observation point is in the plane of the particle orbit, *i.e.*, $\alpha = \pi/2$ the lobe width is given by $\Delta\theta$.

previous page then becomes

$$\frac{dU^{\text{rad}}(\pi/2, 0)}{dt'} = \frac{\mu_0 q'^2 v^2}{16\pi^2 c} \frac{1}{\left(1 - \frac{v}{c}\right)^3} \quad (8.167)$$

which means that an observer near the orbit plane sees a very strong pulse followed, half an orbit period later, by a much weaker pulse.

The two cases represented by equation (8.166) on the facing page and equation (8.167) above are very important results since they can be used to determine the characteristics of the particle motion both in particle accelerators and in astrophysical objects where a direct measurement of particle velocities are impossible.

In the orbit plane ($\alpha = \pi/2$), equation (8.162) on page 143 gives

$$\frac{dU^{\text{rad}}(\pi/2, \varphi)}{dt'} = \frac{\mu_0 q'^2 v^2}{16\pi^2 c} \frac{\left(1 - \frac{v}{c} \cos \varphi\right)^2 - \left(1 - \frac{v^2}{c^2}\right) \sin^2 \varphi}{\left(1 - \frac{v}{c} \cos \varphi\right)^5} \quad (8.168)$$

which vanishes for angles φ_0 which fulfil

$$\cos \varphi_0 = \frac{v}{c} \quad (8.169a)$$

$$\sin \varphi_0 = \sqrt{1 - \frac{v^2}{c^2}} \quad (8.169b)$$

Hence, the angle φ_0 is a measure of the *synchrotron radiation lobe width* $\Delta\theta$; see Figure 8.7. For ultra-relativistic particles, defined by

$$\gamma = \frac{1}{\sqrt{1 - \frac{v^2}{c^2}}} \gg 1, \quad \sqrt{1 - \frac{v^2}{c^2}} \ll 1, \quad (8.170)$$

one can approximate

$$\varphi_0 \approx \sin \varphi_0 = \sqrt{1 - \frac{v^2}{c^2}} = \frac{1}{\gamma} \quad (8.171)$$

Hence, synchrotron radiation from ultra-relativistic charges is characterized by a radiation lobe width which is approximately

$$\Delta\theta \approx \frac{1}{\gamma} \quad (8.172)$$

This angular interval is swept by the charge during the time interval

$$\Delta t' = \frac{\Delta\theta}{\omega_0} \quad (8.173)$$

during which the particle moves a length interval

$$\Delta l = v\Delta t' = v \frac{\Delta\theta}{\omega_0} \quad (8.174)$$

in the direction toward the observer who therefore measures a pulse width of length

$$\begin{aligned} \Delta t &= \Delta t' - \frac{\Delta l}{c} = \Delta t' - \frac{v\Delta t'}{c} = \left(1 - \frac{v}{c}\right) \Delta t' = \left(1 - \frac{v}{c}\right) \frac{\Delta\theta}{\omega_0} \\ &\approx \left(1 - \frac{v}{c}\right) \frac{1}{\gamma\omega_0} = \frac{\left(1 - \frac{v}{c}\right) \left(1 + \frac{v}{c}\right)}{\underbrace{1 + \frac{v}{c}}_{\approx 2}} \frac{1}{\gamma\omega_0} \approx \underbrace{1 - \frac{v^2}{c^2}}_{1/\gamma^2} \frac{1}{2\gamma\omega_0} \\ &= \frac{1}{2\gamma^3} \frac{1}{\omega_0} \end{aligned} \quad (8.175)$$

As a general rule, the spectral width of a pulse of length Δt is $\Delta\omega \lesssim 1/\Delta t$. In the ultra-relativistic synchrotron case one can therefore expect frequency components

up to

$$\omega_{\max} \approx \frac{1}{\Delta t} = 2\gamma^3 \omega_0 \quad (8.176)$$

A spectral analysis of the radiation pulse will exhibit Fourier components $n\omega_0$ from $n = 1$ up to $n \approx 2\gamma^3$.

When N electrons are contributing to the radiation, we can discern between three situations:

1. All electrons are very close to each other so that the individual phase differences are negligible. The power will be multiplied by N^2 relative to a single electron and we talk about *coherent radiation*.
2. The electrons are perfectly evenly distributed in the orbit. This is the case, for instance, for electrons in a circular current in a conductor. In this case the radiation fields cancel completely and no far fields are generated.
3. The electrons are unevenly distributed in the orbit. This happens for an open ring current which is subject to fluctuations of order \sqrt{N} as for all open systems. As a result we get *incoherent radiation*. Examples of this can be found both in earthly laboratories and under cosmic conditions.

Radiation in the general case

We recall that the general expression for the radiation \mathbf{E} field from a moving charge concentration is given by expression (8.92) on page 129. This expression in equation (8.161) on page 143 yields the general formula

$$\frac{dU^{\text{rad}}(\theta)}{dt'} = \frac{\mu_0 q'^2 |\mathbf{x} - \mathbf{x}'|}{16\pi^2 c s^5} \left\{ (\mathbf{x} - \mathbf{x}') \times \left[\left((\mathbf{x} - \mathbf{x}') - \frac{|\mathbf{x} - \mathbf{x}'| \mathbf{v}}{c} \right) \times \dot{\mathbf{v}} \right] \right\}^2 \quad (8.177)$$

Integration over the solid angle Ω gives the totally radiated power as

$$\frac{d\tilde{U}^{\text{rad}}}{dt'} = \frac{\mu_0 q'^2 \dot{v}^2}{6\pi c} \frac{1 - \frac{v^2}{c^2} \sin^2 \theta}{\left(1 - \frac{v^2}{c^2}\right)^3} \quad (8.178)$$

where θ is the angle between \mathbf{v} and $\dot{\mathbf{v}}$.

In the limit $\mathbf{v} \parallel \dot{\mathbf{v}}$, $\sin \theta = 0$, which corresponds to *bremsstrahlung*. For $\mathbf{v} \perp \dot{\mathbf{v}}$, $\sin \theta = 1$, which corresponds to *cyclotron radiation* or *synchrotron radiation*.

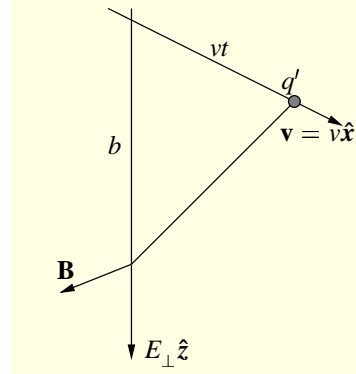


Figure 8.8. The perpendicular field of a charge q' moving with velocity $\mathbf{v} = v\hat{\mathbf{x}}$ is $E_{\perp}\hat{\mathbf{z}}$.

Virtual photons

According to formula (8.103a) on page 132 and Figure 8.8

$$\begin{aligned} E_{\perp} = E_z &= \frac{q'}{4\pi\epsilon_0 s^3} \left(1 - \frac{v^2}{c^2}\right) (\mathbf{x} - \mathbf{x}_0) \cdot \hat{\mathbf{x}}_3 \\ &= \frac{q'}{4\pi\epsilon_0 \gamma^2} \frac{b}{[(vt)^2 + b^2/\gamma^2]^{3/2}} \end{aligned} \quad (8.179)$$

which represents a contracted field, approaching the field of a plane wave. The passage of this field “pulse” corresponds to a frequency distribution of the field energy. Fourier transforming, we obtain

$$E_{\omega,\perp} = \frac{1}{2\pi} \int_{-\infty}^{\infty} E_{\perp}(t) e^{i\omega t} dt = \frac{q}{4\pi^2\epsilon_0 bv} \left[\left(\frac{b\omega}{v\gamma}\right) K_1\left(\frac{b\omega}{v\gamma}\right) \right] \quad (8.180)$$

Here, K_1 is the *Kelvin function* (Bessel function of the second kind with imaginary argument) which behaves in such a way for small and large arguments that

$$E_{\omega,\perp} \sim \frac{q}{4\pi^2\epsilon_0 bv}, \quad b\omega \ll v\gamma \quad (8.181a)$$

$$E_{\omega,\perp} \sim 0, \quad b\omega \gg v\gamma \quad (8.181b)$$

showing that the “pulse” length is of the order $b/v\gamma$.

Due to the equipartition of the field energy into the electric and magnetic fields,

the total field energy can be written

$$U = \varepsilon_0 \int_V E_{\perp}^2 d^3x' = \varepsilon_0 \int_{b_{\min}}^{b_{\max}} \int_{-\infty}^{\infty} E_{\perp}^2 v dt 2\pi b db \quad (8.182)$$

where the volume integration is over the plane perpendicular to \mathbf{v} . With the use of *Parseval's identity* for Fourier transforms, formula (8.15) on page 111, we can rewrite this as

$$\begin{aligned} U &= \int_0^{\infty} U d\omega = 4\pi\varepsilon_0 v \int_{b_{\min}}^{b_{\max}} \int_0^{\infty} |E_{\omega,\perp}|^2 d\omega 2\pi b db \\ &\approx \frac{q^2}{2\pi^2\varepsilon_0 v} \int_0^{\infty} \int_{b_{\min}}^{v\gamma/\omega} \frac{db}{b} d\omega \end{aligned} \quad (8.183)$$

from which we conclude that

$$U_{\omega} \approx \frac{q^2}{2\pi^2\varepsilon_0 v} \ln \left(\frac{v\gamma}{b_{\min}\omega} \right) \quad (8.184)$$

where an explicit value of b_{\min} can be calculated in quantum theory only.

As in the case of bremsstrahlung, it is intriguing to quantise the energy into photons [*cf.* equation (8.155) on page 141]. Then we find that

$$N_{\omega} d\omega \approx \frac{2\alpha}{\pi} \ln \left(\frac{c\gamma}{b_{\min}\omega} \right) \frac{d\omega}{\omega} \quad (8.185)$$

where $\alpha = e^2/(4\pi\varepsilon_0\hbar c) \approx 1/137$ is the *fine structure constant*.

Let us consider the interaction of two electrons, 1 and 2. The result of this interaction is that they change their linear momenta from \mathbf{p}_1 to \mathbf{p}'_1 and \mathbf{p}_2 to \mathbf{p}'_2 , respectively. Heisenberg's uncertainty principle gives $b_{\min} \sim \hbar/|\mathbf{p}_1 - \mathbf{p}'_1|$ so that the number of photons exchanged in the process is of the order

$$N_{\omega} d\omega \approx \frac{2\alpha}{\pi} \ln \left(\frac{c\gamma}{\hbar\omega} |\mathbf{p}_1 - \mathbf{p}'_1| \right) \frac{d\omega}{\omega} \quad (8.186)$$

Since this change in momentum corresponds to a change in energy $\hbar\omega = E_1 - E'_1$ and $E_1 = m_0\gamma c$, we see that

$$N_{\omega} d\omega \approx \frac{2\alpha}{\pi} \ln \left(\frac{E_1}{m_0 c^2} \frac{|c\mathbf{p}_1 - c\mathbf{p}'_1|}{E_1 - E'_1} \right) \frac{d\omega}{\omega} \quad (8.187)$$

a formula which gives a reasonable account of electron- and photon-induced processes.

8.5.5 Radiation from charges moving in matter

When electromagnetic radiation is propagating through matter, new phenomena may appear which are (at least classically) not present in vacuum. As mentioned earlier, one can under certain simplifying assumptions include, to some extent, the influence from matter on the electromagnetic fields by introducing new, derived field quantities \mathbf{D} and \mathbf{H} according to

$$\mathbf{D} = \varepsilon(t, \mathbf{x})\mathbf{E} = \kappa\varepsilon_0\mathbf{E} \quad (8.188)$$

$$\mathbf{B} = \mu(t, \mathbf{x})\mathbf{H} = \kappa_m\mu_0\mathbf{H} \quad (8.189)$$

Expressed in terms of these derived field quantities, the Maxwell equations, often called *macroscopic Maxwell equations*, take the form

$$\nabla \cdot \mathbf{D} = \rho(t, \mathbf{x}) \quad (8.190a)$$

$$\nabla \times \mathbf{E} + \frac{\partial}{\partial t}\mathbf{B} = \mathbf{0} \quad (8.190b)$$

$$\nabla \cdot \mathbf{B} = 0 \quad (8.190c)$$

$$\nabla \times \mathbf{H} - \frac{\partial}{\partial t}\mathbf{D} = \mathbf{j}(t, \mathbf{x}) \quad (8.190d)$$

Assuming for simplicity that the *electric permittivity* ε and the *magnetic permeability* μ , and hence the *relative permittivity* κ and the *relative permeability* κ_m all have fixed values, independent on time and space, for each type of material we consider, we can derive the general *telegrapher's equation* [cf. equation (2.23) on page 27]

$$\frac{\partial^2 \mathbf{E}}{\partial \zeta^2} - \sigma\mu \frac{\partial \mathbf{E}}{\partial t} - \varepsilon\mu \frac{\partial^2 \mathbf{E}}{\partial t^2} = \mathbf{0} \quad (8.191)$$

describing (1D) wave propagation in a material medium.

In chapter 2 we concluded that the existence of a finite conductivity, manifesting itself in a *collisional interaction* between the charge carriers, causes the waves to decay exponentially with time and space. Let us therefore assume that in our medium $\sigma = 0$ so that the wave equation simplifies to

$$\frac{\partial^2 \mathbf{E}}{\partial \zeta^2} - \varepsilon\mu \frac{\partial^2 \mathbf{E}}{\partial t^2} = \mathbf{0} \quad (8.192)$$

If we introduce the *phase velocity* in the medium as

$$v_\phi = \frac{1}{\sqrt{\varepsilon\mu}} = \frac{1}{\sqrt{\kappa\varepsilon_0\kappa_m\mu_0}} = \frac{c}{\sqrt{\kappa\kappa_m}} \quad (8.193)$$

where, according to equation (1.9) on page 5, $c = 1/\sqrt{\varepsilon_0\mu_0}$ is the speed of light, *i.e.*, the phase speed of electromagnetic waves in vacuum, then the general solution

to each component of equation (8.192) on the facing page

$$E_i = f(\zeta - v_\varphi t) + g(\zeta + v_\varphi t), \quad i = 1, 2, 3 \quad (8.194)$$

The ratio of the phase speed in vacuum and in the medium

$$\frac{c}{v_\varphi} = \sqrt{\kappa\kappa_m} = c\sqrt{\varepsilon\mu} \stackrel{\text{def}}{=} n \quad (8.195)$$

is called the *refractive index* of the medium. In general n is a function of both time and space as are the quantities ε , μ , κ , and κ_m themselves. If, in addition, the medium is *anisotropic* or *birefringent*, all these quantities are rank-two tensor fields. Under our simplifying assumptions, in each medium we consider $n = \text{Const}$ for each frequency component of the fields.

Associated with the phase speed of a medium for a wave of a given frequency ω we have a *wave vector*, defined as

$$\mathbf{k} \stackrel{\text{def}}{=} k\hat{\mathbf{k}} = k\hat{\mathbf{v}}_\varphi = \frac{\omega}{v_\varphi} \frac{\mathbf{v}_\varphi}{v_\varphi} \quad (8.196)$$

As in the vacuum case discussed in chapter 2, assuming that \mathbf{E} is time-harmonic, *i.e.*, can be represented by a Fourier component proportional to $\exp\{-i\omega t\}$, the solution of Equation (8.192) can be written

$$\mathbf{E} = \mathbf{E}_0 e^{i(\mathbf{k}\cdot\mathbf{x} - \omega t)} \quad (8.197)$$

where now \mathbf{k} is the wave vector *in the medium* given by equation (8.196) above. With these definitions, the vacuum formula for the associated magnetic field, equation (2.30) on page 28,

$$\mathbf{B} = \sqrt{\varepsilon\mu} \hat{\mathbf{k}} \times \mathbf{E} = \frac{1}{v_\varphi} \hat{\mathbf{k}} \times \mathbf{E} = \frac{1}{\omega} \mathbf{k} \times \mathbf{E} \quad (8.198)$$

is valid also in a material medium (assuming, as mentioned, that n has a fixed constant scalar value). A consequence of a $\kappa \neq 1$ is that the electric field will, in general, have a longitudinal component.

It is important to notice that depending on the electric and magnetic properties of a medium, and, hence, on the value of the refractive index n , the phase speed in the medium can be smaller or larger than the speed of light:

$$v_\varphi = \frac{c}{n} = \frac{\omega}{k} \quad (8.199)$$

where, in the last step, we used equation (8.196) above.

If the medium has a refractive index which, as is usually the case, dependent on frequency ω , we say that the medium is *dispersive*. Because in this case also $\mathbf{k}(\omega)$

and $\omega(\mathbf{k})$, so that the *group velocity*

$$v_g = \frac{\partial \omega}{\partial k} \quad (8.200)$$

has a unique value for each frequency component, and is different from v_ϕ . Except in regions of *anomalous dispersion*, v_ϕ is always smaller than c . In a gas of free charges, such as a *plasma*, the refractive index is given by the expression

$$n^2(\omega) = 1 - \frac{\omega_p^2}{\omega^2} \quad (8.201)$$

where

$$\omega_p^2 = \sum_{\sigma} \frac{N_{\sigma} q_{\sigma}^2}{\epsilon_0 m_{\sigma}} \quad (8.202)$$

is the *plasma frequency*. Here m_{σ} and N_{σ} denote the mass and number density, respectively, of charged particle species σ . In an inhomogeneous plasma, $N_{\sigma} = N_{\sigma}(\mathbf{x})$ so that the refractive index and also the phase and group velocities are space dependent. As can be easily seen, for each given frequency, the phase and group velocities in a plasma are different from each other. If the frequency ω is such that it coincides with ω_p at some point in the medium, then at that point $v_\phi \rightarrow \infty$ while $v_g \rightarrow 0$ and the wave Fourier component at ω is reflected there.

Vavilov-Čerenkov radiation

As we saw in subsection ??, a charge in uniform, rectilinear motion *in vacuum* does not give rise to any radiation; see in particular equation (8.103a) on page 132. Let us now consider a charge in uniform, rectilinear motion *in a medium* with electric properties which are different from those of a (classical) vacuum. Specifically, consider a medium where

$$\epsilon = \text{Const} > \epsilon_0 \quad (8.203a)$$

$$\mu = \mu_0 \quad (8.203b)$$

This implies that in this medium the phase speed is

$$v_\phi = \frac{c}{n} = \frac{1}{\sqrt{\epsilon \mu_0}} < c \quad (8.204)$$

Hence, in this particular medium, the speed of propagation of (the phase planes of) electromagnetic waves is less than the speed of light in vacuum, which we know is

an absolute limit for the motion of anything, including particles. A medium of this kind has the interesting property that particles, entering into the medium at high speeds $|\mathbf{v}|$, which, of course, are below the phase speed *in vacuum*, can experience that the particle speeds are *higher* than the phase speed *in the medium*. This is the basis for the *Vavilov-Čerenkov radiation* that we shall now study.

If we recall the general derivation, in the vacuum case, of the retarded (and advanced) potentials in chapter 3 and the Liénard-Wiechert potentials, equations (8.68) on page 123, we realise that we obtain the latter in the medium by a simple formal replacement $c \rightarrow c/n$ in the expression (8.69) on page 123 for s . Hence, the Liénard-Wiechert potentials in a medium characterized by a refractive index n , are

$$\phi(t, \mathbf{x}) = \frac{1}{4\pi\epsilon_0} \frac{q'}{\left| |\mathbf{x} - \mathbf{x}'| - n \frac{(\mathbf{x} - \mathbf{x}') \cdot \mathbf{v}}{c} \right|} = \frac{1}{4\pi\epsilon_0} \frac{q'}{s} \quad (8.205a)$$

$$\mathbf{A}(t, \mathbf{x}) = \frac{1}{4\pi\epsilon_0 c^2} \frac{q' \mathbf{v}}{\left| |\mathbf{x} - \mathbf{x}'| - n \frac{(\mathbf{x} - \mathbf{x}') \cdot \mathbf{v}}{c} \right|} = \frac{1}{4\pi\epsilon_0 c^2} \frac{q' \mathbf{v}}{s} \quad (8.205b)$$

where now

$$s = \left| |\mathbf{x} - \mathbf{x}'| - n \frac{(\mathbf{x} - \mathbf{x}') \cdot \mathbf{v}}{c} \right| \quad (8.206)$$

The need for the absolute value of the expression for s is obvious in the case when $v/c \geq 1/n$ because then the second term can be larger than the first term; if $v/c \ll 1/n$ we recover the well-known vacuum case but with modified phase speed. We also note that the retarded and advanced times in the medium are [cf. equation (3.34) on page 41]

$$t'_{\text{ret}} = t'_{\text{ret}}(t, |\mathbf{x} - \mathbf{x}'|) = t - \frac{k |\mathbf{x} - \mathbf{x}'|}{\omega} = t - \frac{|\mathbf{x} - \mathbf{x}'| n}{c} \quad (8.207a)$$

$$t'_{\text{adv}} = t'_{\text{adv}}(t, |\mathbf{x} - \mathbf{x}'|) = t + \frac{k |\mathbf{x} - \mathbf{x}'|}{\omega} = t + \frac{|\mathbf{x} - \mathbf{x}'| n}{c} \quad (8.207b)$$

so that the usual time interval $t - t'$ between the time measured at the point of observation and the retarded time *in a medium* becomes

$$t - t' = \frac{|\mathbf{x} - \mathbf{x}'| n}{c} \quad (8.208)$$

For $v/c \geq 1/n$, the retarded distance s , and therefore the denominators in equa-

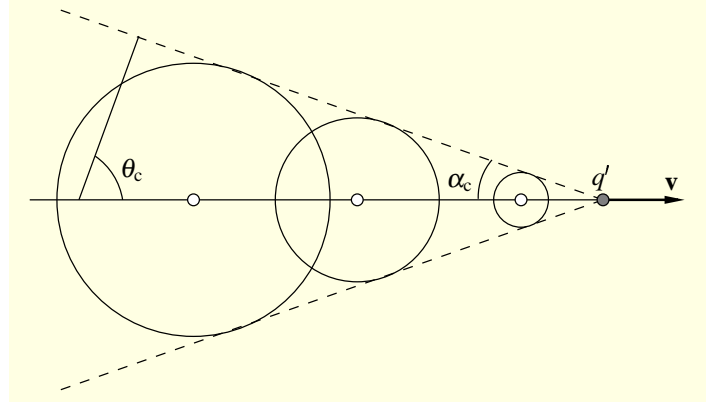


Figure 8.9. Instantaneous picture of the expanding field spheres from a point charge moving with constant speed $v/c > 1/n$ in a medium where $n > 1$. This generates a Vavilov-Čerenkov shock wave in the form of a cone.

tions (8.205) on the preceding page vanish when

$$n(\mathbf{x} - \mathbf{x}') \cdot \frac{\mathbf{v}}{c} = |\mathbf{x} - \mathbf{x}'| \frac{nv}{c} \cos \theta_c = |\mathbf{x} - \mathbf{x}'| \quad (8.209)$$

or, equivalently, when

$$\cos \theta_c = \frac{c}{nv} \quad (8.210)$$

In the direction defined by this angle θ_c , the potentials become singular. During the time interval $t - t'$ given by expression (8.208) on the previous page, the field exists within a sphere of radius $|\mathbf{x} - \mathbf{x}'|$ around the particle while the particle moves a distance

$$l = v(t - t') \quad (8.211)$$

along the direction of \mathbf{v} .

In the direction θ_c where the potentials are singular, all field spheres are tangent to a straight cone with its apex at the instantaneous position of the particle and with the apex half angle α_c defined according to

$$\sin \alpha_c = \cos \theta_c = \frac{c}{nv} \quad (8.212)$$

This cone of potential singularities and field sphere circumferences propagates with

speed c/n in the form of a *shock front*, called *Vavilov-Čerenkov radiation*.¹ The Vavilov-Čerenkov cone is similar in nature to the *Mach cone* in acoustics.

In order to make some quantitative estimates of this radiation, we note that we can describe the motion of each charged particle q' as a current density:

$$\mathbf{j} = q' \mathbf{v} \delta(\mathbf{x}' - \mathbf{v}t') = q' v \delta(x' - vt') \delta(y') \delta(z') \hat{\mathbf{x}}_1 \quad (8.213)$$

which has the trivial Fourier transform

$$\mathbf{j}_\omega = \frac{q'}{2\pi} e^{i\omega x'/v} \delta(y') \delta(z') \hat{\mathbf{x}}_1 \quad (8.214)$$

This Fourier component can be used in the formulae derived for a linear current in subsection 8.3.1 if only we make the replacements

$$\epsilon_0 \rightarrow \epsilon = n^2 \epsilon_0 \quad (8.215a)$$

$$k \rightarrow \frac{n\omega}{c} \quad (8.215b)$$

In this manner, using \mathbf{j}_ω from equation (8.214), the resulting Fourier transforms of the Vavilov-Čerenkov magnetic and electric radiation fields can be calculated from the expressions (8.4) and (8.5) on page 108, respectively.

The total energy content is then obtained from equation (8.15) on page 111 (integrated over a closed sphere at large distances). For a Fourier component one obtains [*cf.* equation (8.18) on page 111]

$$\begin{aligned} U_\omega^{\text{rad}} d\Omega &\approx \frac{1}{4\pi\epsilon_0 n c} \left| \int_V (\mathbf{j}_\omega \times \mathbf{k}) e^{-i\mathbf{k}\cdot\mathbf{x}'} d^3x' \right|^2 d\Omega \\ &= \frac{q'^2 n \omega^2}{16\pi^3 \epsilon_0 c^3} \left| \int_{-\infty}^{\infty} \exp \left[i \left(\frac{\omega x'}{v} - kx' \cos \theta \right) \right] dx' \right|^2 \sin^2 \theta d\Omega \end{aligned} \quad (8.216)$$

where θ is the angle between the direction of motion, $\hat{\mathbf{x}}'_1$, and the direction to the observer, $\hat{\mathbf{k}}$. The integral in (8.216) is singular of a “Dirac delta type.” If we limit

¹The first observation of this radiation was made by P. A. Čerenkov in 1934, who was then a post-graduate student in S. I. Vavilov’s research group at the Lebedev Institute in Moscow. Vavilov wrote a manuscript with the experimental findings, put Čerenkov as the author, and submitted it to *Nature*. In the manuscript, Vavilov explained the results in terms of radioactive particles creating Compton electrons which gave rise to the radiation (which was the correct interpretation), but the paper was rejected. The paper was then sent to *Physical Review* and was, after some controversy with the American editors who claimed the results to be wrong, eventually published in 1937. In the same year, I. E. Tamm and I. M. Frank published the theory for the effect (“the singing electron”). In fact, predictions of a similar effect had been made as early as 1888 by Heaviside, and by Sommerfeld in his 1904 paper “Radiating body moving with velocity of light”. On May 8, 1937, Sommerfeld sent a letter to Tamm via Austria, saying that he was surprised that his old 1904 ideas were now becoming interesting. Tamm, Frank and Čerenkov received the Nobel Prize in 1958 “for the discovery and the interpretation of the Čerenkov effect” [V. L. Ginzburg, *private communication*].

the spatial extent of the motion of the particle to the closed interval $[-X, X]$ on the x' axis we can evaluate the integral to obtain

$$U_{\omega}^{\text{rad}} d\Omega = \frac{q'^2 n \omega^2 \sin^2 \theta}{4\pi^3 \epsilon_0 c^3} \frac{\sin^2 \left[\left(1 - \frac{nv}{c} \cos \theta\right) \frac{X\omega}{v} \right]}{\left[\left(1 - \frac{nv}{c} \cos \theta\right) \frac{\omega}{v} \right]^2} d\Omega \quad (8.217)$$

which has a maximum in the direction θ_c as expected. The magnitude of this maximum grows and its width narrows as $X \rightarrow \infty$. The integration of (8.217) over Ω therefore picks up the main contributions from $\theta \approx \theta_c$. Consequently, we can set $\sin^2 \theta \approx \sin^2 \theta_c$ and the result of the integration is

$$\begin{aligned} \tilde{U}_{\omega}^{\text{rad}} &= 2\pi \int_0^{\pi} U_{\omega}^{\text{rad}} \sin \theta d\theta = 2\pi \int_0^{\pi} U_{\omega}^{\text{rad}} d \underbrace{(-\cos \theta)}_{\xi} \\ &\approx \frac{q'^2 n \omega^2 \sin^2 \theta_c}{2\pi^2 \epsilon_0 c^3} \int_{-1}^1 \frac{\sin^2 \left[\left(1 + \frac{mv\xi}{c}\right) \frac{X\omega}{v} \right]}{\left[\left(1 + \frac{mv\xi}{c}\right) \frac{\omega}{v} \right]^2} d\xi \end{aligned} \quad (8.218)$$

The integral (8.218) is strongly peaked near $\xi = -c/(nv)$, or, equivalently, near $\cos \theta_c = c/(nv)$ so we can extend the integration limits to $\pm\infty$ without introducing too much error. Via yet another variable substitution we can therefore approximate

$$\begin{aligned} \sin^2 \theta_c \int_{-1}^1 \frac{\sin^2 \left[\left(1 + \frac{mv\xi}{c}\right) \frac{X\omega}{v} \right]}{\left[\left(1 + \frac{mv\xi}{c}\right) \frac{\omega}{v} \right]^2} d\xi &\approx \left(1 - \frac{c^2}{n^2 v^2}\right) \frac{cX}{\omega n} \int_{-\infty}^{\infty} \frac{\sin^2 x}{x^2} dx \\ &= \frac{cX\pi}{\omega n} \left(1 - \frac{c^2}{n^2 v^2}\right) \end{aligned} \quad (8.219)$$

leading to the final approximate result for the total energy loss in the frequency interval $(\omega, \omega + d\omega)$

$$\tilde{U}_{\omega}^{\text{rad}} d\omega = \frac{q'^2 X}{2\pi \epsilon_0 c^2} \left(1 - \frac{c^2}{n^2 v^2}\right) \omega d\omega \quad (8.220)$$

As mentioned earlier, the refractive index is usually frequency dependent. Realising this, we find that the radiation energy per frequency unit and *per unit length*

is

$$\frac{\tilde{U}_\omega^{\text{rad}} d\omega}{2X} = \frac{q'^2 \omega}{4\pi\epsilon_0 c^2} \left(1 - \frac{c^2}{n^2(\omega)v^2} \right) d\omega \quad (8.221)$$

This result was derived under the assumption that $v/c > 1/n(\omega)$, *i.e.*, under the condition that the expression inside the parentheses in the right hand side is positive. For all media it is true that $n(\omega) \rightarrow 1$ when $\omega \rightarrow \infty$, so there exist always a highest frequency for which we can obtain Vavilov-Čerenkov radiation from a fast charge in a medium. Our derivation above for a fixed value of n is valid for each individual Fourier component.

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Formulae

F.1 The Electromagnetic Field

F.1.1 Maxwell's equations

$$\nabla \cdot \mathbf{D} = \rho \quad (\text{F.1})$$

$$\nabla \cdot \mathbf{B} = 0 \quad (\text{F.2})$$

$$\nabla \times \mathbf{E} = -\frac{\partial}{\partial t} \mathbf{B} \quad (\text{F.3})$$

$$\nabla \times \mathbf{H} = \mathbf{j} + \frac{\partial}{\partial t} \mathbf{D} \quad (\text{F.4})$$

Constitutive relations

$$\mathbf{D} = \epsilon \mathbf{E} \quad (\text{F.5})$$

$$\mathbf{H} = \frac{\mathbf{B}}{\mu} \quad (\text{F.6})$$

$$\mathbf{j} = \sigma \mathbf{E} \quad (\text{F.7})$$

$$\mathbf{P} = \epsilon_0 \chi \mathbf{E} \quad (\text{F.8})$$

F.1.2 Fields and potentials

Vector and scalar potentials

$$\mathbf{B} = \nabla \times \mathbf{A} \quad (\text{F.9})$$

$$\mathbf{E} = -\nabla\phi - \frac{\partial}{\partial t}\mathbf{A} \quad (\text{F.10})$$

Lorentz' gauge condition in vacuum

$$\nabla \cdot \mathbf{A} + \frac{1}{c^2} \frac{\partial}{\partial t} \phi = 0 \quad (\text{F.11})$$

F.1.3 Force and energy

Poynting's vector

$$\mathbf{S} = \mathbf{E} \times \mathbf{H} \quad (\text{F.12})$$

Maxwell's stress tensor

$$T_{ij} = E_i D_j + H_i B_j - \frac{1}{2} \delta_{ij} (E_k D_k + H_k B_k) \quad (\text{F.13})$$

F.2 Electromagnetic Radiation

F.2.1 Relationship between the field vectors in a plane wave

$$\mathbf{B} = \frac{\hat{\mathbf{k}} \times \mathbf{E}}{c} \quad (\text{F.14})$$

F.2.2 The far fields from an extended source distribution

$$\mathbf{B}_\omega^{\text{rad}}(\mathbf{x}) = \frac{-i\mu_0}{4\pi} \frac{e^{ik|\mathbf{x}|}}{|\mathbf{x}|} \int_{V'} d^3x' e^{-i\mathbf{k}\cdot\mathbf{x}'} \mathbf{j}_\omega \times \mathbf{k} \quad (\text{F.15})$$

$$\mathbf{E}_\omega^{\text{rad}}(\mathbf{x}) = \frac{i}{4\pi\epsilon_0 c} \frac{e^{ik|\mathbf{x}|}}{|\mathbf{x}|} \hat{\mathbf{x}} \times \int_{V'} d^3x' e^{-i\mathbf{k}\cdot\mathbf{x}'} \mathbf{j}_\omega \times \mathbf{k} \quad (\text{F.16})$$

F.2.3 The far fields from an electric dipole

$$\mathbf{B}_\omega^{\text{rad}}(\mathbf{x}) = -\frac{\omega\mu_0}{4\pi} \frac{e^{ik|\mathbf{x}|}}{|\mathbf{x}|} \mathbf{p}_\omega \times \mathbf{k} \quad (\text{F.17})$$

$$\mathbf{E}_\omega^{\text{rad}}(\mathbf{x}) = -\frac{1}{4\pi\epsilon_0} \frac{e^{ik|\mathbf{x}|}}{|\mathbf{x}|} (\mathbf{p}_\omega \times \mathbf{k}) \times \mathbf{k} \quad (\text{F.18})$$

F.2.4 The far fields from a magnetic dipole

$$\mathbf{B}_\omega^{\text{rad}}(\mathbf{x}) = -\frac{\mu_0}{4\pi} \frac{e^{ik|\mathbf{x}|}}{|\mathbf{x}|} (\mathbf{m}_\omega \times \mathbf{k}) \times \mathbf{k} \quad (\text{F.19})$$

$$\mathbf{E}_\omega^{\text{rad}}(\mathbf{x}) = \frac{k}{4\pi\epsilon_0 c} \frac{e^{ik|\mathbf{x}|}}{|\mathbf{x}|} \mathbf{m}_\omega \times \mathbf{k} \quad (\text{F.20})$$

F.2.5 The far fields from an electric quadrupole

$$\mathbf{B}_\omega^{\text{rad}}(\mathbf{x}) = \frac{i\mu_0\omega}{8\pi} \frac{e^{ik|\mathbf{x}|}}{|\mathbf{x}|} (\mathbf{k} \cdot \mathbf{Q}_\omega) \times \mathbf{k} \quad (\text{F.21})$$

$$\mathbf{E}_\omega^{\text{rad}}(\mathbf{x}) = \frac{i}{8\pi\epsilon_0} \frac{e^{ik|\mathbf{x}|}}{|\mathbf{x}|} [(\mathbf{k} \cdot \mathbf{Q}_\omega) \times \mathbf{k}] \times \mathbf{k} \quad (\text{F.22})$$

F.2.6 The fields from a point charge in arbitrary motion

$$\mathbf{E}(t, \mathbf{x}) = \frac{q}{4\pi\epsilon_0 s^3} \left[\mathbf{R}_v \left(1 - \frac{v^2}{c^2} \right) + (\mathbf{x} - \mathbf{x}') \times \frac{\mathbf{R}_v \times \dot{\mathbf{v}}}{c^2} \right] \quad (\text{F.23})$$

$$\mathbf{B}(t, \mathbf{x}) = (\mathbf{x} - \mathbf{x}') \times \frac{\mathbf{E}(t, \mathbf{x})}{c|\mathbf{x} - \mathbf{x}'|} \quad (\text{F.24})$$

$$s = |\mathbf{x} - \mathbf{x}'| - (\mathbf{x} - \mathbf{x}') \cdot \frac{\mathbf{v}}{c} \quad (\text{F.25})$$

$$\mathbf{R}_v = (\mathbf{x} - \mathbf{x}') - |\mathbf{x} - \mathbf{x}'| \frac{\mathbf{v}}{c} \quad (\text{F.26})$$

$$\left(\frac{\partial t'}{\partial t} \right)_{\mathbf{x}} = \frac{|\mathbf{x} - \mathbf{x}'|}{s} \quad (\text{F.27})$$

F.2.7 The fields from a point charge in uniform motion

$$\mathbf{E}(t, \mathbf{x}) = \frac{q}{4\pi\epsilon_0 s^3} \left(1 - \frac{v^2}{c^2}\right) \mathbf{R}_0 \quad (\text{F.28})$$

$$\mathbf{B}(t, \mathbf{x}) = \frac{\mathbf{v} \times \mathbf{E}(t, \mathbf{x})}{c^2} \quad (\text{F.29})$$

$$s = \sqrt{|\mathbf{R}_0|^2 - \left(\frac{\mathbf{R}_0 \times \mathbf{v}}{c}\right)^2} \quad (\text{F.30})$$

$$\mathbf{R}_0 = \mathbf{x} - \mathbf{x}_0 \quad (\text{F.31})$$

F.3 Special Relativity

F.3.1 Metric tensor

$$g_{\mu\nu} = \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & -1 & 0 & 0 \\ 0 & 0 & -1 & 0 \\ 0 & 0 & 0 & -1 \end{pmatrix} \quad (\text{F.32})$$

F.3.2 Covariant and contravariant four-vectors

$$v_\mu = g_{\mu\nu} v^\nu \quad (\text{F.33})$$

F.3.3 Lorentz transformation of a four-vector

$$x'^\mu = \Lambda^\mu_\nu x^\nu \quad (\text{F.34})$$

$$\Lambda^\mu_\nu = \begin{pmatrix} \gamma & -\gamma\beta & 0 & 0 \\ -\gamma\beta & \gamma & 0 & 0 \\ 0 & 0 & 1 & 0 \\ 0 & 0 & 0 & 1 \end{pmatrix} \quad (\text{F.35})$$

$$\gamma = \frac{1}{\sqrt{1 - \beta^2}} \quad (\text{F.36})$$

$$\beta = \frac{v}{c} \quad (\text{F.37})$$

F.3.4 Invariant line element

$$ds = c \frac{dt}{\gamma} = c d\tau \quad (\text{F.38})$$

F.3.5 Four-velocity

$$u^\mu = \frac{dx^\mu}{ds} = \left(\gamma, \gamma \frac{\mathbf{v}}{c} \right) \quad (\text{F.39})$$

F.3.6 Four-momentum

$$p^\mu = m_0 c^2 u^\mu = (E, c\mathbf{p}) \quad (\text{F.40})$$

F.3.7 Four-current density

$$j^\mu = \rho_0 u^\mu = \left(\rho, \rho \frac{\mathbf{v}}{c} \right) \quad (\text{F.41})$$

F.3.8 Four-potential

$$A^\mu = (\phi, c\mathbf{A}) \quad (\text{F.42})$$

F.3.9 Field tensor

$$F^{\mu\nu} = \frac{\partial A^\nu}{\partial x_\mu} - \frac{\partial A^\mu}{\partial x_\nu} = \begin{pmatrix} 0 & -E_x & -E_y & -E_z \\ E_x & 0 & -cB_z & cB_y \\ E_y & cB_z & 0 & -cB_x \\ E_z & -cB_y & cB_x & 0 \end{pmatrix} \quad (\text{F.43})$$

F.4 Vector Relations

Let \mathbf{x} be the radius vector (coordinate vector), from the origin to the point $(x_1, x_2, x_3) \equiv (x, y, z)$ and let $|\mathbf{x}|$ denote the magnitude (“length”) of \mathbf{x} . Let further $\alpha(\mathbf{x}), \beta(\mathbf{x}), \dots$ be arbitrary scalar fields and $\mathbf{a}(\mathbf{x}), \mathbf{b}(\mathbf{x}), \mathbf{c}(\mathbf{x}), \mathbf{d}(\mathbf{x}), \dots$ arbitrary vector fields.

The differential vector operator ∇ is in Cartesian coordinates given by

$$\nabla \equiv \sum_{i=1}^3 \stackrel{\text{def}}{\equiv} \hat{\mathbf{x}}_i \frac{\partial}{\partial x_i} \stackrel{\text{def}}{\equiv} \boldsymbol{\partial} \quad (\text{F.44})$$

where $\hat{\mathbf{x}}_i$, $i = 1, 2, 3$ is the i th unit vector and $\hat{\mathbf{x}}_1 \equiv \hat{\mathbf{x}}$, $\hat{\mathbf{x}}_2 \equiv \hat{\mathbf{y}}$, and $\hat{\mathbf{x}}_3 \equiv \hat{\mathbf{z}}$. In

component (tensor) notation ∇ can be written

$$\nabla_i = \partial_i = \left(\frac{\partial}{\partial x_1}, \frac{\partial}{\partial x_2}, \frac{\partial}{\partial x_3} \right) = \left(\frac{\partial}{\partial x}, \frac{\partial}{\partial y}, \frac{\partial}{\partial z} \right) \quad (\text{F.45})$$

F.4.1 Spherical polar coordinates

Base vectors

$$\hat{\mathbf{r}} = \sin \theta \cos \varphi \hat{\mathbf{x}} + \sin \theta \sin \varphi \hat{\mathbf{y}} + \cos \theta \hat{\mathbf{z}} \quad (\text{F.46})$$

$$\hat{\boldsymbol{\theta}} = \cos \theta \cos \varphi \hat{\mathbf{x}} + \cos \theta \sin \varphi \hat{\mathbf{y}} - \sin \theta \hat{\mathbf{z}} \quad (\text{F.47})$$

$$\hat{\boldsymbol{\phi}} = -\sin \varphi \hat{\mathbf{x}} + \cos \varphi \hat{\mathbf{y}} \quad (\text{F.48})$$

Directed line element

$$dx \hat{\mathbf{x}} = d\mathbf{l} = dr \hat{\mathbf{r}} + r d\theta \hat{\boldsymbol{\theta}} + r^2 \sin \theta d\varphi \hat{\boldsymbol{\phi}} \quad (\text{F.49})$$

Solid angle element

$$d\Omega = \sin \theta d\theta d\varphi \quad (\text{F.50})$$

Directed area element

$$d^2x \hat{\mathbf{n}} = d\mathbf{S} = dS \hat{\mathbf{r}} = r^2 d\Omega \hat{\mathbf{r}} \quad (\text{F.51})$$

Volume element

$$d^3x = dV = dr dS = r^2 dr d\Omega \quad (\text{F.52})$$

F.4.2 Vector formulae

General relations

$$\mathbf{a} \cdot \mathbf{b} = \mathbf{b} \cdot \mathbf{a} = \delta_{ij} a_i b_j = ab \cos \theta \quad (\text{F.53})$$

$$\mathbf{a} \times \mathbf{b} = -\mathbf{b} \times \mathbf{a} = \varepsilon_{ijk} a_j b_k \hat{\mathbf{x}}_i \quad (\text{F.54})$$

$$\mathbf{a} \cdot (\mathbf{b} \times \mathbf{c}) = (\mathbf{a} \times \mathbf{b}) \cdot \mathbf{c} \quad (\text{F.55})$$

$$\mathbf{a} \times (\mathbf{b} \times \mathbf{c}) = \mathbf{b}(\mathbf{a} \cdot \mathbf{c}) - \mathbf{c}(\mathbf{a} \cdot \mathbf{b}) \quad (\text{F.56})$$

$$\mathbf{a} \times (\mathbf{b} \times \mathbf{c}) + \mathbf{b} \times (\mathbf{c} \times \mathbf{a}) + \mathbf{c} \times (\mathbf{a} \times \mathbf{b}) = \mathbf{0} \quad (\text{F.57})$$

$$(\mathbf{a} \times \mathbf{b}) \cdot (\mathbf{c} \times \mathbf{d}) = \mathbf{a} \cdot [\mathbf{b} \times (\mathbf{c} \times \mathbf{d})] = (\mathbf{a} \cdot \mathbf{c})(\mathbf{b} \cdot \mathbf{d}) - (\mathbf{a} \cdot \mathbf{d})(\mathbf{b} \cdot \mathbf{c}) \quad (\text{F.58})$$

$$(\mathbf{a} \times \mathbf{b}) \times (\mathbf{c} \times \mathbf{d}) = (\mathbf{a} \times \mathbf{b} \cdot \mathbf{d})\mathbf{c} - (\mathbf{a} \times \mathbf{b} \cdot \mathbf{c})\mathbf{d} \quad (\text{F.59})$$

$$\nabla(\alpha\beta) = \alpha\nabla\beta + \beta\nabla\alpha \quad (\text{F.60})$$

$$\nabla \cdot (\alpha\mathbf{a}) = \mathbf{a} \cdot \nabla\alpha + \alpha\nabla \cdot \mathbf{a} \quad (\text{F.61})$$

$$\nabla \times (\alpha\mathbf{a}) = \alpha\nabla \times \mathbf{a} - \mathbf{a} \times \nabla\alpha \quad (\text{F.62})$$

$$\nabla \cdot (\mathbf{a} \times \mathbf{b}) = \mathbf{b} \cdot (\nabla \times \mathbf{a}) - \mathbf{a} \cdot (\nabla \times \mathbf{b}) \quad (\text{F.63})$$

$$\nabla \times (\mathbf{a} \times \mathbf{b}) = \mathbf{a}(\nabla \cdot \mathbf{b}) - \mathbf{b}(\nabla \cdot \mathbf{a}) + (\mathbf{b} \cdot \nabla)\mathbf{a} - (\mathbf{a} \cdot \nabla)\mathbf{b} \quad (\text{F.64})$$

$$\nabla (\mathbf{a} \cdot \mathbf{b}) = \mathbf{a} \times (\nabla \times \mathbf{b}) + \mathbf{b} \times (\nabla \times \mathbf{a}) + (\mathbf{b} \cdot \nabla)\mathbf{a} + (\mathbf{a} \cdot \nabla)\mathbf{b} \quad (\text{F.65})$$

$$\nabla \cdot \nabla \alpha = \nabla^2 \alpha \quad (\text{F.66})$$

$$\nabla \times \nabla \alpha = \mathbf{0} \quad (\text{F.67})$$

$$\nabla \cdot (\nabla \times \mathbf{a}) = 0 \quad (\text{F.68})$$

$$\nabla \times (\nabla \times \mathbf{a}) = \nabla(\nabla \cdot \mathbf{a}) - \nabla^2 \mathbf{a} \quad (\text{F.69})$$

Special relations

In the following \mathbf{k} is an arbitrary *constant* vector.

$$\nabla \cdot \mathbf{x} = 3 \quad (\text{F.70})$$

$$\nabla \times \mathbf{x} = \mathbf{0} \quad (\text{F.71})$$

$$\nabla |\mathbf{x}| = \frac{\mathbf{x}}{|\mathbf{x}|} \quad (\text{F.72})$$

$$\nabla \left(\frac{1}{|\mathbf{x}|} \right) = -\frac{\mathbf{x}}{|\mathbf{x}|^3} \quad (\text{F.73})$$

$$\nabla \cdot \left(\frac{\mathbf{x}}{|\mathbf{x}|^3} \right) = -\nabla^2 \left(\frac{1}{|\mathbf{x}|} \right) = 4\pi \delta(\mathbf{x}) \quad (\text{F.74})$$

$$\nabla \left(\frac{\mathbf{k}}{|\mathbf{x}|} \right) = \mathbf{k} \cdot \left[\nabla \left(\frac{1}{|\mathbf{x}|} \right) \right] = -\frac{\mathbf{k} \cdot \mathbf{x}}{|\mathbf{x}|^3} \quad (\text{F.75})$$

$$\nabla \times \left[\mathbf{k} \times \left(\frac{\mathbf{x}}{|\mathbf{x}|^3} \right) \right] = -\nabla \left(\frac{\mathbf{k} \cdot \mathbf{x}}{|\mathbf{x}|^3} \right) \text{ if } |\mathbf{x}| \neq 0 \quad (\text{F.76})$$

$$\nabla^2 \left(\frac{\mathbf{k}}{|\mathbf{x}|} \right) = \mathbf{k} \nabla^2 \left(\frac{1}{|\mathbf{x}|} \right) = -4\pi \mathbf{k} \delta(\mathbf{x}) \quad (\text{F.77})$$

$$\nabla \times (\mathbf{k} \times \mathbf{a}) = \mathbf{k}(\nabla \cdot \mathbf{a}) + \mathbf{k} \times (\nabla \times \mathbf{a}) - \nabla(\mathbf{k} \cdot \mathbf{a}) \quad (\text{F.78})$$

Integral relations

Let $V(S)$ be the volume bounded by the closed surface $S(V)$. Denote the 3-dimensional volume element by $d^3x (\equiv dV)$ and the surface element, directed along the outward pointing surface normal unit vector $\hat{\mathbf{n}}$, by $d\mathbf{S} (\equiv d^2x \hat{\mathbf{n}})$.

$$\int_V (\nabla \cdot \mathbf{a}) d^3x = \oint_S d\mathbf{S} \cdot \mathbf{a} \quad (\text{F.79})$$

$$\int_V (\nabla \alpha) d^3x = \oint_S d\mathbf{S} \alpha \quad (\text{F.80})$$

$$\int_V (\nabla \times \mathbf{a}) d^3x = \oint_S d\mathbf{S} \times \mathbf{a} \quad (\text{F.81})$$

If $S(C)$ is an open surface bounded by the contour $C(S)$, whose line element is $d\mathbf{l}$, then

$$\oint_C \alpha d\mathbf{l} = \int_S d\mathbf{S} \times \nabla \alpha \quad (\text{F.82})$$

$$\oint_C \mathbf{a} \cdot d\mathbf{l} = \int_S d\mathbf{S} \cdot (\nabla \times \mathbf{a}) \quad (\text{F.83})$$

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Mathematical Methods

M.1 Scalars, Vectors and Tensors

Every physical observable can be described by a geometric object. We will describe the observables in classical electrodynamics mathematically in terms of scalars, pseudoscalars, vectors, pseudovectors, tensors or pseudotensors and will not exploit differential forms to any significant degree.

A *scalar* describes a scalar quantity which may or may not be constant in time and/or space. A *vector* describes some kind of physical motion due tovection and a *tensor* describes the motion or deformation due to some form of tension. However, generalisations to more abstract notions of these quantities are commonplace. The difference between a scalar, vector and tensor and a *pseudoscalar*, *pseudovector* and a *pseudotensor* is that the latter behave differently under such coordinate transformations which cannot be reduced to pure rotations.

Throughout we adopt the convention that Latin indices i, j, k, l, \dots run over the range 1, 2, 3 to denote vector or tensor components in the real Euclidean three-dimensional (3D) configuration space \mathbb{R}^3 , and Greek indices $\mu, \nu, \kappa, \lambda, \dots$, which are used in four-dimensional (4D) space, run over the range 0, 1, 2, 3.

M.1.1 Vectors

Radius vector

Any vector can be represented mathematically in several different ways. One suitable representation is in terms of an ordered N -tuple, or *row vector*, of the coordinates x_N where N is the dimensionality of the space under consideration. The most basic vector is *radius vector* which is the vector from the origin to the point of interest. Its N -tuple representation simply enumerates the coordinates which

describe this point. In this sense, the radius vector from the origin to a point is synonymous with the coordinates of the point itself.

In the 3D space \mathbb{R}^3 , we have $N = 3$ and the radius vector can be represented by the triplet (x_1, x_2, x_3) of coordinates x_i , $i = 1, 2, 3$. The coordinates x_i are scalar quantities which describe the position along the unit base vectors $\hat{\mathbf{x}}_i$ which span \mathbb{R}^3 . Therefore a representation of the radius vector in \mathbb{R}^3 is

$$\mathbf{x} = \sum_{i=1}^3 \hat{\mathbf{x}}_i x_i \stackrel{\text{def}}{\equiv} \hat{\mathbf{x}}_i x_i \quad (\text{M.1})$$

where we have introduced *Einstein's summation convention* (E Σ) which states that a repeated index in a term implies summation over the range of the index in question. Whenever possible and convenient we shall in the following always assume E Σ and suppress explicit summation in our formulae. Typographically, we represent a 3D vector by a boldface letter or symbol in a Roman font.

Alternatively, we may describe the radius vector in *component notation* as follows:

$$x_i \stackrel{\text{def}}{\equiv} (x_1, x_2, x_3) \equiv (x, y, z) \quad (\text{M.2})$$

This component notation is particularly useful in 4D space where we can represent the radius vector either in its *contravariant component form*

$$x^\mu \stackrel{\text{def}}{\equiv} (x^0, x^1, x^2, x^3) \quad (\text{M.3})$$

or its *covariant component form*

$$x_\mu \stackrel{\text{def}}{\equiv} (x_0, x_1, x_2, x_3) \quad (\text{M.4})$$

The relation between the covariant and contravariant forms is determined by the *metric tensor* (also known as the *fundamental tensor*) whose actual form is dictated by the physics. The dual representation of vectors in contravariant and covariant forms is most convenient when we work in a non-Euclidean vector space with an indefinite *metric*. An example is *Lorentz space* \mathbb{L}^4 which is a 4D *Riemannian space*. \mathbb{L}^4 is often utilised to formulate the special theory of relativity.

We note that for a change of coordinates $x^\mu \rightarrow x'^\mu = x'^\mu(x^0, x^1, x^2, x^3)$, due to a transformation from a system Σ to another system Σ' , the differential radius vector

dx^μ transforms as

$$dx'^\mu = \frac{\partial x'^\mu}{\partial x^\nu} dx^\nu \quad (\text{M.5})$$

which follows trivially from the rules of differentiation of x'^μ considered as functions of four variables x^ν .

M.1.2 Fields

A *field* is a physical entity which depends on one or more continuous parameters. Such a parameter can be viewed as a “continuous index” which enumerates the “coordinates” of the field. In particular, in a field which depends on the usual radius vector \mathbf{x} of \mathbb{R}^3 , each point in this space can be considered as one degree of freedom so that a field is a representation of a physical entity which has an infinite number of degrees of freedom.

Scalar fields

We denote an arbitrary *scalar field* in \mathbb{R}^3 by

$$\alpha(\mathbf{x}) = \alpha(x_1, x_2, x_3) \stackrel{\text{def}}{\equiv} \alpha(x_i) \quad (\text{M.6})$$

This field describes how the scalar quantity α varies continuously in 3D \mathbb{R}^3 space.

In 4D, a *four-scalar* field is denoted

$$\alpha(x^0, x^1, x^2, x^3) \stackrel{\text{def}}{\equiv} \alpha(x^\mu) \quad (\text{M.7})$$

which indicates that the four-scalar α depends on all four coordinates spanning this space. Since a four-scalar has the same value at a given point regardless of coordinate system, it is also called an *invariant*.

Analogous to the transformation rule, equation (M.5), for the differential dx^μ , the transformation rule for the differential operator $\partial/\partial x^\mu$ under a transformation

$x^\mu \rightarrow x'^\mu$ becomes

$$\frac{\partial}{\partial x'^\mu} = \frac{\partial x^\nu}{\partial x'^\mu} \frac{\partial}{\partial x^\nu} \quad (\text{M.8})$$

which, again, follows trivially from the rules of differentiation.

Vector fields

We can represent an arbitrary vector field $\mathbf{a}(\mathbf{x})$ in \mathbb{R}^3 as follows:

$$\mathbf{a}(\mathbf{x}) = \hat{\mathbf{x}}_i a_i(\mathbf{x}) \quad (\text{M.9})$$

In component notation this same vector can be represented as

$$a_i(\mathbf{x}) = (a_1(\mathbf{x}), a_2(\mathbf{x}), a_3(\mathbf{x})) = a_i(x_j) \quad (\text{M.10})$$

In 4D, an arbitrary *four-vector* field in contravariant component form can be represented as

$$a^\mu(x^\nu) = (a^0(x^\nu), a^1(x^\nu), a^2(x^\nu), a^3(x^\nu)) \quad (\text{M.11})$$

or, in *covariant* component form, as

$$a_\mu(x^\nu) = (a_0(x^\nu), a_1(x^\nu), a_2(x^\nu), a_3(x^\nu)) \quad (\text{M.12})$$

where x^ν is the radius four-vector. Again, the relation between a^μ and a_μ is determined by the metric of the physical 4D system under consideration.

Whether an arbitrary N -tuple fulfils the requirement of being an (N -dimensional) contravariant vector or not, depends on its transformation properties during a change of coordinates. For instance, in 4D an assemblage $y^\mu = (y^0, y^1, y^2, y^3)$ constitutes a *contravariant four-vector* (or the contravariant components of a four-vector) if and only if, during a transformation from a system Σ with coordinates x^μ to a system Σ' with coordinates x'^μ , it transforms to the new system according to the rule

$$y'^\mu = \frac{\partial x'^\mu}{\partial x^\nu} y^\nu \quad (\text{M.13})$$

i.e., in the same way as the differential coordinate element dx^μ transforms according to equation (M.5) on the previous page.

The analogous requirement for a *covariant four-vector* is that it transforms, dur-

ing the change from Σ to Σ' , according to the rule

$$y'_\mu = \frac{\partial x^\nu}{\partial x'^\mu} y_\nu \quad (\text{M.14})$$

i.e., in the same way as the differential operator $\partial/\partial x^\mu$ transforms according to equation (M.8) on the facing page.

Tensor fields

We denote an arbitrary *tensor field* in \mathbb{R}^3 by $\mathbf{A}(\mathbf{x})$. This tensor field can be represented in a number of ways, for instance in the following *matrix form*:

$$\mathbf{A}(\mathbf{x}) = \begin{pmatrix} A_{11}(\mathbf{x}) & A_{12}(\mathbf{x}) & A_{13}(\mathbf{x}) \\ A_{21}(\mathbf{x}) & A_{22}(\mathbf{x}) & A_{23}(\mathbf{x}) \\ A_{31}(\mathbf{x}) & A_{32}(\mathbf{x}) & A_{33}(\mathbf{x}) \end{pmatrix} \stackrel{\text{def}}{\equiv} A_{ij}(x_k) \quad (\text{M.15})$$

where, in the last member, we have again used the more compact component notation. Strictly speaking, the tensor field described here is a tensor of *rank two*.

A particularly simple rank-two tensor in \mathbb{R}^3 is the 3D *Kronecker delta* symbol δ_{ij} , with the following properties:

$$\delta_{ij} = \begin{cases} 0 & \text{if } i \neq j \\ 1 & \text{if } i = j \end{cases} \quad (\text{M.16})$$

The 3D Kronecker delta has the following matrix representation

$$(\delta_{ij}) = \begin{pmatrix} 1 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & 1 \end{pmatrix} \quad (\text{M.17})$$

Another common and useful tensor is the fully antisymmetric tensor of rank 3, also known as the *Levi-Civita tensor*

$$\varepsilon_{ijk} = \begin{cases} 1 & \text{if } i, j, k \text{ is an even permutation of } 1,2,3 \\ 0 & \text{if at least two of } i, j, k \text{ are equal} \\ -1 & \text{if } i, j, k \text{ is an odd permutation of } 1,2,3 \end{cases} \quad (\text{M.18})$$

with the following further property

$$\varepsilon_{ijk}\varepsilon_{ilm} = \delta_{jl}\delta_{km} - \delta_{jm}\delta_{kl} \quad (\text{M.19})$$

In fact, tensors may have any rank n . In this picture a scalar is considered to be a tensor of rank $n = 0$ and a vector a tensor of rank $n = 1$. Consequently, the

notation where a vector (tensor) is represented in its component form is called the *tensor notation*. A tensor of rank $n = 2$ may be represented by a two-dimensional array or matrix whereas higher rank tensors are best represented in their component forms (tensor notation).

In 4D, we have three forms of *four-tensor fields* of rank n . We speak of

- a *contravariant four-tensor field*, denoted $A^{\mu_1\mu_2\cdots\mu_n}(x^V)$,
- a *covariant four-tensor field*, denoted $A_{\mu_1\mu_2\cdots\mu_n}(x^V)$,
- a *mixed four-tensor field*, denoted $A^{\mu_1\mu_2\cdots\mu_k}_{\mu_{k+1}\cdots\mu_n}(x^V)$.

The 4D *metric tensor (fundamental tensor)* mentioned above is a particularly important four-tensor of rank 2. In covariant component form we shall denote it $g_{\mu\nu}$. This metric tensor determines the relation between an arbitrary contravariant four-vector a^μ and its covariant counterpart a_μ according to the following rule:

$$a_\mu(x^\kappa) \stackrel{\text{def}}{=} g_{\mu\nu}a^\nu(x^\kappa) \quad (\text{M.20})$$

This rule is often called *lowering of index*. The *raising of index* analogue of the index lowering rule is:

$$a^\mu(x^\kappa) \stackrel{\text{def}}{=} g^{\mu\nu}a_\nu(x^\kappa) \quad (\text{M.21})$$

More generally, the following lowering and raising rules hold for arbitrary rank n mixed tensor fields:

$$g_{\mu_k\nu_k}A^{\nu_1\nu_2\cdots\nu_{k-1}\nu_k}_{\nu_{k+1}\nu_{k+2}\cdots\nu_n}(x^\kappa) = A^{\nu_1\nu_2\cdots\nu_{k-1}}_{\mu_k\nu_{k+1}\cdots\nu_n}(x^\kappa) \quad (\text{M.22})$$

$$g^{\mu_k\nu_k}A^{\nu_1\nu_2\cdots\nu_{k-1}\nu_k}_{\nu_{k+1}\nu_{k+2}\cdots\nu_n}(x^\kappa) = A^{\nu_1\nu_2\cdots\nu_{k-1}\mu_k}_{\nu_{k+1}\nu_{k+2}\cdots\nu_n}(x^\kappa) \quad (\text{M.23})$$

Successive lowering and raising of more than one index is achieved by a repeated application of this rule. For example, a dual application of the lowering operation on a rank 2 tensor in contravariant form yields

$$A_{\mu\nu} = g_{\mu\kappa}g_{\lambda\nu}A^{\kappa\lambda} \quad (\text{M.24})$$

i.e., the same rank 2 tensor in covariant form. This operation is also known as a *tensor contraction*.

▷ TENSORS IN 3D SPACE

EXAMPLE M.1

Consider the tetrahedron-like volume element V indicated in Figure M.1 on the next page of a solid, fluid, or gaseous body, whose atomistic structure is irrelevant for the present analysis. Let $d\mathbf{S} = d^2x \hat{\mathbf{n}}$ in Figure M.1 on the following page be a directed surface element of this volume element and let the vector $\mathbf{T}_{\hat{\mathbf{n}}} d^2x$ be the force that matter, lying on the side of d^2x toward which the unit normal vector $\hat{\mathbf{n}}$ points, acts on matter which lies on the opposite side of d^2x . This force concept is meaningful only if the forces are short-range enough that they can be assumed to act only in the surface proper. According to Newton's third law, this surface force fulfils

$$\mathbf{T}_{-\hat{\mathbf{n}}} = -\mathbf{T}_{\hat{\mathbf{n}}} \quad (\text{M.25})$$

Using (M.25) and Newton's second law, we find that the matter of mass m , which at a given instant is located in V obeys the equation of motion

$$\mathbf{T}_{\hat{\mathbf{n}}} d^2x - \mathbf{T}_{\hat{\mathbf{x}}_1} d^2x - \mathbf{T}_{\hat{\mathbf{x}}_2} d^2x - \mathbf{T}_{\hat{\mathbf{x}}_3} d^2x + \mathbf{F}_{\text{ext}} = m\mathbf{a} \quad (\text{M.26})$$

where \mathbf{F}_{ext} is the external force and \mathbf{a} is the acceleration of the volume element. In other words

$$\mathbf{T}_{\hat{\mathbf{n}}} = n_1 \mathbf{T}_{\hat{\mathbf{x}}_1} + n_2 \mathbf{T}_{\hat{\mathbf{x}}_2} + n_3 \mathbf{T}_{\hat{\mathbf{x}}_3} + \frac{m}{d^2x} \left(\mathbf{a} - \frac{\mathbf{F}_{\text{ext}}}{m} \right) \quad (\text{M.27})$$

Since both \mathbf{a} and $\mathbf{F}_{\text{ext}}/m$ remain finite whereas $m/d^2x \rightarrow 0$ as $V \rightarrow 0$, one finds that in this limit

$$\mathbf{T}_{\hat{\mathbf{n}}} = \sum_{i=1}^3 n_i \mathbf{T}_{\hat{\mathbf{x}}_i} \equiv n_i \mathbf{T}_{\hat{\mathbf{x}}_i} \quad (\text{M.28})$$

From the above derivation it is clear that equation (M.28) above is valid not only in equilibrium but also when the matter in V is in motion.

Introducing the notation

$$T_{ij} = \left(\mathbf{T}_{\hat{\mathbf{x}}_i} \right)_j \quad (\text{M.29})$$

for the j th component of the vector $\mathbf{T}_{\hat{\mathbf{x}}_i}$, we can write equation (M.28) in component form as follows

$$T_{\hat{\mathbf{n}}j} = \left(\mathbf{T}_{\hat{\mathbf{n}}} \right)_j = \sum_{i=1}^3 n_i T_{ij} \equiv n_i T_{ij} \quad (\text{M.30})$$

Using equation (M.30) above, we find that the component of the vector $\mathbf{T}_{\hat{\mathbf{n}}}$ in the direction of an arbitrary unit vector $\hat{\mathbf{m}}$ is

$$\begin{aligned} T_{\hat{\mathbf{n}}\hat{\mathbf{m}}} &= \mathbf{T}_{\hat{\mathbf{n}}} \cdot \hat{\mathbf{m}} \\ &= \sum_{j=1}^3 T_{\hat{\mathbf{n}}j} m_j = \sum_{j=1}^3 \left(\sum_{i=1}^3 n_i T_{ij} \right) m_j \equiv n_i T_{ij} m_j = \hat{\mathbf{n}} \cdot \mathbf{T} \cdot \hat{\mathbf{m}} \end{aligned} \quad (\text{M.31})$$

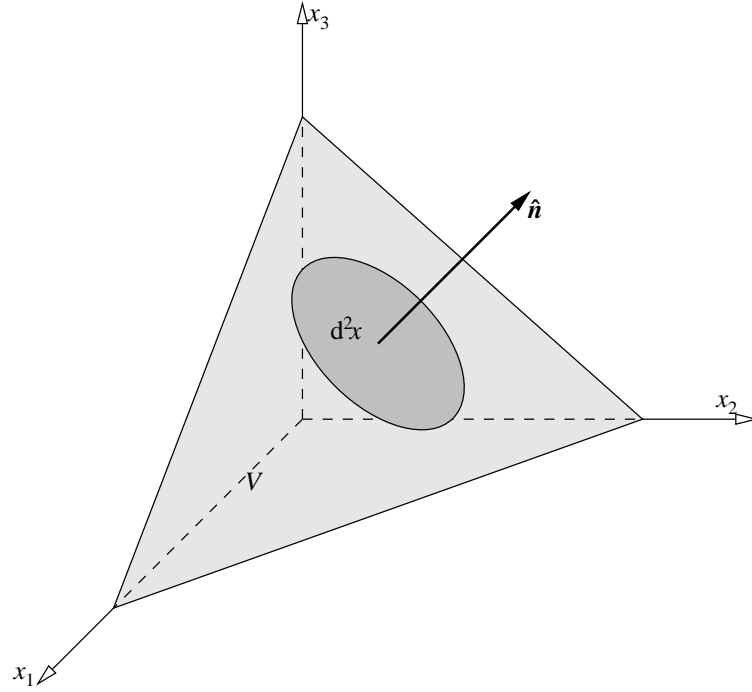


Figure M.1. Tetrahedron-like volume element V containing matter.

Hence, the j th component of the vector $\mathbf{T}_{\hat{x}_i}$, here denoted T_{ij} , can be interpreted as the ij th component of a tensor \mathbf{T} . Note that $T_{\hat{n}\hat{m}}$ is independent of the particular coordinate system used in the derivation.

We shall now show how one can use the momentum law (force equation) to derive the equation of motion for an arbitrary element of mass in the body. To this end we consider a part V of the body. If the external force density (force per unit volume) is denoted by \mathbf{f} and the velocity for a mass element dm is denoted by \mathbf{v} , we obtain

$$\frac{d}{dt} \int_V \mathbf{v} dm = \int_V \mathbf{f} d^3x + \int_S \mathbf{T}_{\hat{n}} d^2x \quad (\text{M.32})$$

The j th component of this equation can be written

$$\int_V \frac{d}{dt} v_j dm = \int_V f_j d^3x + \int_S T_{\hat{n}j} d^2x = \int_V f_j d^3x + \int_S n_i T_{ij} d^2x \quad (\text{M.33})$$

where, in the last step, equation (M.30) on the preceding page was used. Setting $dm = \rho d^3x$ and using the divergence theorem on the last term, we can rewrite the result as

$$\int_V \rho \frac{d}{dt} v_j d^3x = \int_V f_j d^3x + \int_V \frac{\partial T_{ij}}{\partial x_i} d^3x \quad (\text{M.34})$$

Since this formula is valid for any arbitrary volume, we must require that

$$\rho \frac{d}{dt} v_j - f_j - \frac{\partial T_{ij}}{\partial x_i} = 0 \quad (\text{M.35})$$

or, equivalently

$$\rho \frac{\partial v_j}{\partial t} + \rho \mathbf{v} \cdot \nabla v_j - f_j - \frac{\partial T_{ij}}{\partial x_i} = 0 \quad (\text{M.36})$$

Note that $\partial v_j / \partial t$ is the rate of change with time of the velocity component v_j at a *fixed* point $\mathbf{x} = (x_1, x_2, x_3)$.

END OF EXAMPLE M.1 ◀

M.1.3 Vector algebra

Scalar product

The *scalar product* (*dot product*, *inner product*) of two arbitrary 3D vectors \mathbf{a} and \mathbf{b} in ordinary \mathbb{R}^3 space is the scalar number

$$\mathbf{a} \cdot \mathbf{b} = \hat{\mathbf{x}}_i a_i \cdot \hat{\mathbf{x}}_j b_j = \hat{\mathbf{x}}_i \cdot \hat{\mathbf{x}}_j a_i b_j = \delta_{ij} a_i b_j = a_i b_i \quad (\text{M.37})$$

where we used the fact that the scalar product $\hat{\mathbf{x}}_i \cdot \hat{\mathbf{x}}_j$ is a representation of the Kronecker delta δ_{ij} defined in equation (M.16) on page 177. In Russian literature, the scalar product is often denoted (\mathbf{ab}) .

In 4D space we define the scalar product of two arbitrary four-vectors a^μ and b^μ in the following way

$$a_\mu b^\mu = g_{\nu\mu} a^\nu b^\mu = a^\nu b_\nu = g^{\mu\nu} a_\mu b_\nu \quad (\text{M.38})$$

where we made use of the index lowering and raising rules (M.20) and (M.21). The result is a four-scalar, *i.e.*, an invariant which is independent on in which inertial system it is measured.

The *quadratic differential form*

$$ds^2 = g_{\mu\nu} dx^\nu dx^\mu = dx_\mu dx^\mu \quad (\text{M.39})$$

i.e., the scalar product of the differential radius four-vector with itself, is an invariant called the *metric*. It is also the square of the *line element* ds which is the distance between neighbouring points with coordinates x^μ and $x^\mu + dx^\mu$.

EXAMPLE M.2

▷ INNER PRODUCT IN COMPLEX VECTOR SPACE

A 3D *complex vector* \mathbf{A} is a vector in \mathbb{C}^3 (or, if we like, in \mathbb{R}^6), expressed in terms of two real vectors \mathbf{a}_R and \mathbf{a}_I in \mathbb{R}^3 in the following way

$$\mathbf{A} \stackrel{\text{def}}{\equiv} \mathbf{a}_R + i\mathbf{a}_I = a_R \hat{\mathbf{a}}_R + ia_I \hat{\mathbf{a}}_I \stackrel{\text{def}}{\equiv} A \hat{\mathbf{A}} \in \mathbb{C}^3 \quad (\text{M.40})$$

The inner product of \mathbf{A} with itself may be defined as

$$\mathbf{A}^2 \stackrel{\text{def}}{\equiv} \mathbf{A} \cdot \mathbf{A} = a_R^2 - a_I^2 + 2i\mathbf{a}_R \cdot \mathbf{a}_I \stackrel{\text{def}}{\equiv} A^2 \in \mathbb{C} \quad (\text{M.41})$$

from which we find that

$$A = \sqrt{a_R^2 - a_I^2 + 2i\mathbf{a}_R \cdot \mathbf{a}_I} \in \mathbb{C} \quad (\text{M.42})$$

Using this in equation (M.40), we see that we can interpret this so that the complex unit vector is

$$\hat{\mathbf{A}} = \frac{\mathbf{A}}{A} = \frac{a_R}{\sqrt{a_R^2 - a_I^2 + 2i\mathbf{a}_R \cdot \mathbf{a}_I}} \hat{\mathbf{a}}_R + i \frac{a_I}{\sqrt{a_R^2 - a_I^2 + 2i\mathbf{a}_R \cdot \mathbf{a}_I}} \hat{\mathbf{a}}_I \in \mathbb{C}^3 \quad (\text{M.43})$$

On the other hand, the usual definition of the scalar product of a vector in complex vector space with itself is

$$|\mathbf{A}|^2 \stackrel{\text{def}}{\equiv} \mathbf{A} \cdot \mathbf{A}^* = a_R^2 + a_I^2 = |A|^2 \in \mathbb{R} \quad (\text{M.44})$$

—END OF EXAMPLE M.2◁

EXAMPLE M.3

▷ SCALAR PRODUCT, NORM AND METRIC IN LORENTZ SPACE

In \mathbb{L}^4 the metric tensor attains a simple form [see equation (5.7) on page 58 for an example] and, hence, the scalar product in equation (M.38) on the preceding page can be evaluated almost trivially and becomes

$$a_\mu b^\mu = (a_0, -\mathbf{a}) \cdot (b^0, \mathbf{b}) = a_0 b^0 - \mathbf{a} \cdot \mathbf{b} \quad (\text{M.45})$$

The important scalar product of the \mathbb{L}^4 radius four-vector with itself becomes

$$\begin{aligned} x_\mu x^\mu &= (x_0, -\mathbf{x}) \cdot (x^0, \mathbf{x}) = (ct, -\mathbf{x}) \cdot (ct, \mathbf{x}) \\ &= (ct)^2 - (x^1)^2 - (x^2)^2 - (x^3)^2 = s^2 \end{aligned} \quad (\text{M.46})$$

which is the indefinite, real *norm* of \mathbb{L}^4 . The \mathbb{L}^4 metric is the quadratic differential form

$$ds^2 = dx_\mu dx^\mu = c^2(dt)^2 - (dx^1)^2 - (dx^2)^2 - (dx^3)^2 \quad (\text{M.47})$$

—END OF EXAMPLE M.3◁

▷ METRIC IN GENERAL RELATIVITY

EXAMPLE M.4

In the general theory of relativity, several important problems are treated in a 4D spherical polar coordinate system where the radius four-vector can be given as $x^\mu = (ct, r, \theta, \phi)$ and the metric tensor is

$$g_{\mu\nu} = \begin{pmatrix} e^\kappa & 0 & 0 & 0 \\ 0 & e^{-\lambda} & 0 & 0 \\ 0 & 0 & -r^2 & 0 \\ 0 & 0 & 0 & -r^2 \sin^2 \theta \end{pmatrix} \quad (\text{M.48})$$

where $\kappa = \kappa(ct, r, \theta, \phi)$ and $\lambda = \lambda(ct, r, \theta, \phi)$. In such a space, the metric takes the form

$$ds^2 = c^2 e^\kappa (dt)^2 - e^\lambda (dr)^2 - r^2 (d\theta)^2 - r^2 \sin^2 \theta (d\phi)^2 \quad (\text{M.49})$$

In general relativity the metric tensor is not given *a priori* but is determined by the *Einstein equations*.

—END OF EXAMPLE M.4◁

Dyadic product

A tensor $\mathbf{A}(\mathbf{x})$ can sometimes be represented in the *dyadic form* $\mathbf{A}(\mathbf{x}) \equiv \mathbf{a}(\mathbf{x})\mathbf{b}(\mathbf{x})$. The dyadic notation with two juxtaposed vectors \mathbf{a} and \mathbf{b} is interpreted as an *outer product* and this dyad is operated on by another vector \mathbf{c} from the right and from the left with a scalar (inner) product in the following two ways:

$$\mathbf{A} \cdot \mathbf{c} \stackrel{\text{def}}{\equiv} \mathbf{a}\mathbf{b} \cdot \mathbf{c} \stackrel{\text{def}}{\equiv} \mathbf{a}(\mathbf{b} \cdot \mathbf{c}) \quad (\text{M.50a})$$

$$\mathbf{c} \cdot \mathbf{A} \stackrel{\text{def}}{\equiv} \mathbf{c} \cdot \mathbf{a}\mathbf{b} \stackrel{\text{def}}{\equiv} (\mathbf{c} \cdot \mathbf{a})\mathbf{b} \quad (\text{M.50b})$$

thus producing new vectors, proportional to \mathbf{a} and \mathbf{b} . In mathematics, a dyadic product is often called *tensor product* and is frequently denoted $\mathbf{a} \otimes \mathbf{b}$.

Vector product

The *vector product* or *cross product* of two arbitrary 3D vectors \mathbf{a} and \mathbf{b} in ordinary \mathbb{R}^3 space is the vector

$$\mathbf{c} = \mathbf{a} \times \mathbf{b} = \varepsilon_{ijk} a_j b_k \hat{\mathbf{x}}_i \quad (\text{M.51})$$

Here ε_{ijk} is the Levi-Civita tensor defined in equation (M.18) on page 177. Sometimes the vector product of \mathbf{a} and \mathbf{b} is denoted $\mathbf{a} \wedge \mathbf{b}$ or, particularly in the Russian literature, $[\mathbf{a}\mathbf{b}]$.

A spatial reversal of the coordinate system $(x'_1, x'_2, x'_3) = (-x_1, -x_2, -x_3)$ changes sign of the components of the vectors \mathbf{a} and \mathbf{b} so that in the new coordinate system

$\mathbf{a}' = -\mathbf{a}$ and $\mathbf{b}' = -\mathbf{b}$, which is to say that the direction of an ordinary vector is not dependent on the choice of directions of the coordinate axes. On the other hand, as is seen from equation (M.51) on the preceding page, the cross product vector \mathbf{c} does not change sign. Therefore \mathbf{a} (or \mathbf{b}) is an example of a “true” vector, or *polar vector*, whereas \mathbf{c} is an example of an *axial vector*, or *pseudovector*.

A prototype for a pseudovector is the angular momentum vector and hence the attribute “axial.” Pseudovectors transform as ordinary vectors under translations and proper rotations, but reverse their sign relative to ordinary vectors for any coordinate change involving reflection. Tensors (of any rank) which transform analogously to pseudovectors are called *pseudotensors*. Scalars are tensors of rank zero, and zero-rank pseudotensors are therefore also called *pseudoscalars*, an example being the pseudoscalar $\hat{\mathbf{x}}_i \cdot (\hat{\mathbf{x}}_j \times \hat{\mathbf{x}}_k)$. This triple product is a representation of the ijk component of the Levi-Civita tensor ε_{ijk} which is a rank three pseudotensor.

M.1.4 Vector analysis

The del operator

In \mathbb{R}^3 the *del operator* is a *differential vector operator*, denoted in *Gibbs’ notation* by ∇ and defined as

$$\nabla \stackrel{\text{def}}{=} \hat{\mathbf{x}}_i \frac{\partial}{\partial x_i} \stackrel{\text{def}}{=} \boldsymbol{\partial} \quad (\text{M.52})$$

where $\hat{\mathbf{x}}_i$ is the i th unit vector in a Cartesian coordinate system. Since the operator in itself has vectorial properties, we denote it with a boldface nabla. In “component” notation we can write

$$\partial_i = \left(\frac{\partial}{\partial x_1}, \frac{\partial}{\partial x_2}, \frac{\partial}{\partial x_3} \right) \quad (\text{M.53})$$

In 4D, the contravariant component representation of the *four-del operator* is defined by

$$\partial^\mu = \left(\frac{\partial}{\partial x_0}, \frac{\partial}{\partial x_1}, \frac{\partial}{\partial x_2}, \frac{\partial}{\partial x_3} \right) \quad (\text{M.54})$$

whereas the covariant four-del operator is

$$\partial_\mu = \left(\frac{\partial}{\partial x^0}, \frac{\partial}{\partial x^1}, \frac{\partial}{\partial x^2}, \frac{\partial}{\partial x^3} \right) \quad (\text{M.55})$$

We can use this four-del operator to express the transformation properties (M.13)

and (M.14) on page 177 as

$$y'^{\mu} = (\partial_{\nu} x'^{\mu}) y^{\nu} \quad (\text{M.56})$$

and

$$y'_{\mu} = (\partial'_{\mu} x^{\nu}) y_{\nu} \quad (\text{M.57})$$

respectively.

▷ THE FOUR-DEL OPERATOR IN LORENTZ SPACE

EXAMPLE M.5

In \mathbb{L}^4 the contravariant form of the four-del operator can be represented as

$$\partial^{\mu} = \left(\frac{1}{c} \frac{\partial}{\partial t}, -\boldsymbol{\partial} \right) = \left(\frac{1}{c} \frac{\partial}{\partial t}, -\boldsymbol{\nabla} \right) \quad (\text{M.58})$$

and the covariant form as

$$\partial_{\mu} = \left(\frac{1}{c} \frac{\partial}{\partial t}, \boldsymbol{\partial} \right) = \left(\frac{1}{c} \frac{\partial}{\partial t}, \boldsymbol{\nabla} \right) \quad (\text{M.59})$$

Taking the scalar product of these two, one obtains

$$\partial^{\mu} \partial_{\mu} = \frac{1}{c^2} \frac{\partial^2}{\partial t^2} - \nabla^2 = \square^2 \quad (\text{M.60})$$

which is the *d'Alembert operator*, sometimes denoted \square , and sometimes defined with an opposite sign convention.

END OF EXAMPLE M.5◁

With the help of the del operator we can define the gradient, divergence and curl of a tensor (in the generalised sense).

The gradient

The *gradient* of an \mathbb{R}^3 scalar field $\alpha(\mathbf{x})$, denoted $\boldsymbol{\nabla} \alpha(x)$, is an \mathbb{R}^3 vector field $\mathbf{a}(\mathbf{x})$:

$$\boldsymbol{\nabla} \alpha(\mathbf{x}) = \boldsymbol{\partial} \alpha(\mathbf{x}) = \hat{\mathbf{x}}_i \partial_i \alpha(\mathbf{x}) = \mathbf{a}(\mathbf{x}) \quad (\text{M.61})$$

From this we see that the boldface notation for the nabla and del operators is very handy as it elucidates the 3D vectorial property of the gradient.

In 4D, the *four-gradient* is a covariant vector, formed as a derivative of a four-scalar field $\alpha(x^{\mu})$, with the following component form:

$$\partial_{\mu} \alpha(x^{\nu}) = \frac{\partial \alpha(x^{\nu})}{\partial x^{\mu}} \quad (\text{M.62})$$

EXAMPLE M.6

▷ GRADIENTS OF SCALAR FUNCTIONS OF RELATIVE DISTANCES IN 3D

Very often electrodynamic quantities are dependent on the relative distance in \mathbb{R}^3 between two vectors \mathbf{x} and \mathbf{x}' , *i.e.*, on $|\mathbf{x} - \mathbf{x}'|$. In analogy with equation (M.52) on page 184, we can define the “primed” del operator in the following way:

$$\nabla' = \hat{x}_i \frac{\partial}{\partial x'_i} = \partial' \quad (\text{M.63})$$

Using this, the “unprimed” version, equation (M.52) on page 184, and elementary rules of differentiation, we obtain the following two very useful results:

$$\begin{aligned} \nabla (|\mathbf{x} - \mathbf{x}'|) &= \hat{x}_i \frac{\partial |\mathbf{x} - \mathbf{x}'|}{\partial x_i} \\ &= \frac{\mathbf{x} - \mathbf{x}'}{|\mathbf{x} - \mathbf{x}'|} \\ &= -\hat{x}_i \frac{\partial |\mathbf{x} - \mathbf{x}'|}{\partial x'_i} \\ &= -\nabla' (|\mathbf{x} - \mathbf{x}'|) \end{aligned} \quad (\text{M.64})$$

$$\nabla \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right) = -\frac{\mathbf{x} - \mathbf{x}'}{|\mathbf{x} - \mathbf{x}'|^3} = -\nabla' \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right) \quad (\text{M.65})$$

◁ END OF EXAMPLE M.6

The divergence

We define the 3D *divergence* of a vector field in \mathbb{R}^3 as

$$\nabla \cdot \mathbf{a}(\mathbf{x}) = \partial \cdot \hat{x}_j a_j(\mathbf{x}) = \delta_{ij} \partial_i a_j(\mathbf{x}) = \partial_i a_i(\mathbf{x}) = \frac{\partial a_i(\mathbf{x})}{\partial x_i} = \alpha(\mathbf{x}) \quad (\text{M.66})$$

which, as indicated by the notation $\alpha(\mathbf{x})$, is a *scalar* field in \mathbb{R}^3 . We may think of the divergence as a scalar product between a vectorial operator and a vector. As is the case for any scalar product, the result of a divergence operation is a scalar. Again we see that the boldface notation for the 3D del operator is very convenient.

The *four-divergence* of a four-vector a^μ is the following four-scalar:

$$\partial_\mu a^\mu(x^\nu) = \partial^\mu a_\mu(x^\nu) = \frac{\partial a^\mu(x^\nu)}{\partial x^\mu} \quad (\text{M.67})$$

▷ DIVERGENCE IN 3D

EXAMPLE M.7

For an arbitrary \mathbb{R}^3 vector field $\mathbf{a}(\mathbf{x}')$, the following relation holds:

$$\nabla' \cdot \left(\frac{\mathbf{a}(\mathbf{x}')}{|\mathbf{x} - \mathbf{x}'|} \right) = \frac{\nabla' \cdot \mathbf{a}(\mathbf{x}')}{|\mathbf{x} - \mathbf{x}'|} + \mathbf{a}(\mathbf{x}') \cdot \nabla' \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right) \quad (\text{M.68})$$

which demonstrates how the “primed” divergence, defined in terms of the “primed” del operator in equation (M.63) on the facing page, works.

—END OF EXAMPLE M.7◁

The Laplacian

The 3D *Laplace operator* or *Laplacian* can be described as the divergence of the gradient operator:

$$\nabla^2 = \Delta = \nabla \cdot \nabla = \frac{\partial}{\partial x_i} \hat{\mathbf{x}}_i \cdot \hat{\mathbf{x}}_j \frac{\partial}{\partial x_j} = \delta_{ij} \partial_i \partial_j = \partial_i^2 = \frac{\partial^2}{\partial x_i^2} \equiv \sum_{i=1}^3 \frac{\partial^2}{\partial x_i^2} \quad (\text{M.69})$$

The symbol ∇^2 is sometimes read *del squared*. If, for a scalar field $\alpha(\mathbf{x})$, $\nabla^2 \alpha < 0$ at some point in 3D space, it is a sign of *concentration* of α at that point.

▷ THE LAPLACIAN AND THE DIRAC DELTA

EXAMPLE M.8

A very useful formula in 3D \mathbb{R}^3 is

$$\nabla \cdot \nabla \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right) = \nabla^2 \left(\frac{1}{|\mathbf{x} - \mathbf{x}'|} \right) = -4\pi \delta(\mathbf{x} - \mathbf{x}') \quad (\text{M.70})$$

where $\delta(\mathbf{x} - \mathbf{x}')$ is the 3D *Dirac delta* “function.”

—END OF EXAMPLE M.8◁

The curl

In \mathbb{R}^3 the *curl* of a vector field $\mathbf{a}(\mathbf{x})$, denoted $\nabla \times \mathbf{a}(\mathbf{x})$, is another \mathbb{R}^3 vector field $\mathbf{b}(\mathbf{x})$ which can be defined in the following way:

$$\nabla \times \mathbf{a}(\mathbf{x}) = \varepsilon_{ijk} \hat{\mathbf{x}}_i \partial_j a_k(\mathbf{x}) = \varepsilon_{ijk} \hat{\mathbf{x}}_i \frac{\partial a_k(\mathbf{x})}{\partial x_j} = \mathbf{b}(\mathbf{x}) \quad (\text{M.71})$$

where use was made of the Levi-Civita tensor, introduced in equation (M.18) on page 177.

The covariant 4D generalisation of the curl of a four-vector field $a^\mu(x^\nu)$ is the antisymmetric four-tensor field

$$G_{\mu\nu}(x^\kappa) = \partial_\mu a_\nu(x^\kappa) - \partial_\nu a_\mu(x^\kappa) = -G_{\nu\mu}(x^\kappa) \quad (\text{M.72})$$

A vector with vanishing curl is said to be *irrotational*.

EXAMPLE M.9

▷ THE CURL OF A GRADIENT

Using the definition of the \mathbb{R}^3 curl, equation (M.71) on the preceding page, and the gradient, equation (M.61) on page 185, we see that

$$\nabla \times [\nabla \alpha(\mathbf{x})] = \varepsilon_{ijk} \hat{\mathbf{x}}_i \partial_j \partial_k \alpha(\mathbf{x}) \quad (\text{M.73})$$

which, due to the assumed well-behavedness of $\alpha(\mathbf{x})$, vanishes:

$$\begin{aligned} \varepsilon_{ijk} \hat{\mathbf{x}}_i \partial_j \partial_k \alpha(\mathbf{x}) &= \varepsilon_{ijk} \frac{\partial}{\partial x_j} \frac{\partial}{\partial x_k} \alpha(\mathbf{x}) \hat{\mathbf{x}}_i \\ &= \left(\frac{\partial^2}{\partial x_2 \partial x_3} - \frac{\partial^2}{\partial x_3 \partial x_2} \right) \alpha(\mathbf{x}) \hat{\mathbf{x}}_1 \\ &\quad + \left(\frac{\partial^2}{\partial x_3 \partial x_1} - \frac{\partial^2}{\partial x_1 \partial x_3} \right) \alpha(\mathbf{x}) \hat{\mathbf{x}}_2 \\ &\quad + \left(\frac{\partial^2}{\partial x_1 \partial x_2} - \frac{\partial^2}{\partial x_2 \partial x_1} \right) \alpha(\mathbf{x}) \hat{\mathbf{x}}_3 \\ &\equiv \mathbf{0} \end{aligned} \quad (\text{M.74})$$

We thus find that

$$\nabla \times [\nabla \alpha(\mathbf{x})] \equiv \mathbf{0} \quad (\text{M.75})$$

for any arbitrary, well-behaved \mathbb{R}^3 scalar field $\alpha(\mathbf{x})$.

In 4D we note that for any well-behaved four-scalar field $\alpha(x^\kappa)$

$$(\partial_\mu \partial_\nu - \partial_\nu \partial_\mu) \alpha(x^\kappa) \equiv 0 \quad (\text{M.76})$$

so that the four-curl of a four-gradient vanishes just as does a curl of a gradient in \mathbb{R}^3 .

Hence, a gradient is always irrotational.

END OF EXAMPLE M.9◁

EXAMPLE M.10

▷ THE DIVERGENCE OF A CURL

With the use of the definitions of the divergence (M.66) and the curl, equation (M.71) on the preceding page, we find that

$$\nabla \cdot [\nabla \times \mathbf{a}(\mathbf{x})] = \partial_i [\nabla \times \mathbf{a}(\mathbf{x})]_i = \varepsilon_{ijk} \partial_i \partial_j a_k(\mathbf{x}) \quad (\text{M.77})$$

Using the definition for the Levi-Civita symbol, defined by equation (M.18) on page 177,

we find that, due to the assumed well-behavedness of $\mathbf{a}(x)$,

$$\begin{aligned} \partial_i \varepsilon_{ijk} \partial_j a_k(\mathbf{x}) &= \frac{\partial}{\partial x_i} \varepsilon_{ijk} \frac{\partial}{\partial x_j} a_k \\ &= \left(\frac{\partial^2}{\partial x_2 \partial x_3} - \frac{\partial^2}{\partial x_3 \partial x_2} \right) a_1(\mathbf{x}) \\ &\quad + \left(\frac{\partial^2}{\partial x_3 \partial x_1} - \frac{\partial^2}{\partial x_1 \partial x_3} \right) a_2(\mathbf{x}) \\ &\quad + \left(\frac{\partial^2}{\partial x_1 \partial x_2} - \frac{\partial^2}{\partial x_2 \partial x_1} \right) a_3(\mathbf{x}) \\ &\equiv 0 \end{aligned} \tag{M.78}$$

i.e., that

$$\nabla \cdot [\nabla \times \mathbf{a}(\mathbf{x})] \equiv 0 \tag{M.79}$$

for any arbitrary, well-behaved \mathbb{R}^3 vector field $\mathbf{a}(\mathbf{x})$.

In 4D, the four-divergence of the four-curl is *not* zero, for

$$\partial^\nu G_{\mu\nu} = \partial^\mu \partial_\nu a^\nu(x^\kappa) - \square^2 a^\mu(x^\kappa) \neq 0 \tag{M.80}$$

—END OF EXAMPLE M.10◁

Numerous vector algebra and vector analysis formulae are given in chapter F. Those which are not found there can often be easily derived by using the component forms of the vectors and tensors, together with the Kronecker and Levi-Civita tensors and their generalisations to higher ranks. A short but very useful reference in this respect is the article by A. Evett [?].

M.2 Analytical Mechanics

M.2.1 Lagrange's equations

As is well known from elementary analytical mechanics, the *Lagrange function* or *Lagrangian* L is given by

$$L(q_i, \dot{q}_i, t) = L\left(q_i, \frac{dq_i}{dt}, t\right) = T - V \tag{M.81}$$

where q_i is the *generalised coordinate*, T the *kinetic energy* and V the *potential energy* of a mechanical system, The Lagrangian satisfies the *Lagrange equations*

$$\frac{\partial}{\partial t} \left(\frac{\partial L}{\partial \dot{q}_i} \right) - \frac{\partial L}{\partial q_i} = 0 \tag{M.82}$$

We define the to the generalised coordinate q_i *canonically conjugate momentum*

p_i according to

$$p_i = \frac{\partial L}{\partial \dot{q}_i} \quad (\text{M.83})$$

and note from equation (M.82) on the preceding page that

$$\frac{\partial L}{\partial q_i} = \dot{p}_i \quad (\text{M.84})$$

M.2.2 Hamilton's equations

From L , the *Hamiltonian (Hamilton function) H* can be defined via the *Legendre transformation*

$$H(p_i, q_i, t) = p_i \dot{q}_i - L(q_i, \dot{q}_i, t) \quad (\text{M.85})$$

After differentiating the left and right hand sides of this definition and setting them equal we obtain

$$\frac{\partial H}{\partial p_i} dp_i + \frac{\partial H}{\partial q_i} dq_i + \frac{\partial H}{\partial t} dt = \dot{q}_i dp_i + p_i d\dot{q}_i - \frac{\partial L}{\partial q_i} dq_i - \frac{\partial L}{\partial \dot{q}_i} d\dot{q}_i - \frac{\partial L}{\partial t} dt \quad (\text{M.86})$$

According to the definition of p_i , equation (M.83) above, the second and fourth terms on the right hand side cancel. Furthermore, noting that according to equation (M.84) the third term on the right hand side of equation (M.86) above is equal to $-\dot{p}_i dq_i$ and identifying terms, we obtain the *Hamilton equations*:

$$\frac{\partial H}{\partial p_i} = \dot{q}_i = \frac{dq_i}{dt} \quad (\text{M.87a})$$

$$\frac{\partial H}{\partial q_i} = -\dot{p}_i = -\frac{dp_i}{dt} \quad (\text{M.87b})$$

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