

PHYS415

GENERAL RELATIVITY

Lecture Notes

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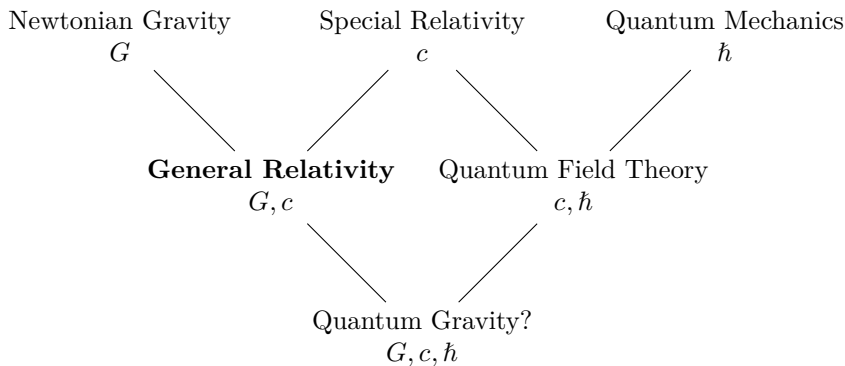
0 Physical Overview of General Relativity

IN GENERAL RELATIVITY, gravity is no longer a “force”...

Space tells matter how to move;

Matter tells space how to curve.

— Misner, Thorne and Wheeler



In Newtonian gravity, escape velocity is given by $\frac{1}{2}mv^2 = \frac{GMm}{R}$. Gravity is significant if $\frac{v^2}{2} \sim \frac{GM}{R}$. Special relativity is significant when $v^2 \sim c^2$. Hence, general relativity is significant when $\frac{c^2}{2} \sim \frac{GM}{R} \iff R \sim \frac{2GM}{c^2}$; this is the *Schwarzschild radius*.

When escape velocity \sim light velocity, the existence of a black hole is implied.

I Review of Special Relativity

— Background —

See (Schutz, 2009, ch 1) and (Doughty, 2018, ch 5, 12, 13).

Assumptions of Special Relativity:

1. The world is described by a 4-dimensional continuum¹, **spacetime**, or **Minkowski space** \mathcal{M}^4 , which is the set of all **events** x^μ ,

$$\vec{x}.$$

— Notation —

Greek indices, μ, ν run over *spacetime* index values; $\{0, 1, 2, 3\}$.

Latin indices (mid-alphabet), i, j, k run over *spatial* index values; $\{1, 2, 3\}$.

\underline{x} is a 4-vector (twiddle); \vec{x} is a 3-vector (under bar).

2. There exist **inertial frames**; namely, frames in which the measured values of time t and position x^i of events result in *linear* equations of motion for *free* particles.

Postulates of Special Relativity:

1. **PRINCIPLE OF RELATIVITY**: The laws of physics are invariant under transformations $x^\nu \rightarrow x^{\bar{\mu}}(x^\nu)$ from one inertial frame to another (and such transformations form a group).
2. **CONSTANCY OF THE SPEED OF LIGHT**: There exists an invariant upper bound on all velocities

$$\left| \frac{d\vec{x}}{dt} \right| \leq \left| \frac{d\vec{x}}{dt} \right|_{\max} = c \quad (\text{speed of light})$$

and this value is the same in all inertial frames.

¹Mathematically, a *Lorentzian manifold*...

For photons,

$$c = \left| \frac{d\vec{x}}{dt} \right|_{\text{frame } K} = \left| \frac{d\vec{x}'}{dt} \right|_{\text{frame } K'}. \quad (\text{I.1})$$

Rewriting (I.1) we see that, for photons,

$$(dx)^2 + (dy)^2 + (dz)^2 - c^2(dt)^2 = (dx')^2 + (dy')^2 + (dz')^2 - c^2(dt')^2 = 0.$$

$\text{frame } K$ $\text{frame } K'$

This suggests the definition of the **spacetime interval** between any neighbouring events x^μ and $x^\mu + dx^\mu$ as

$$\begin{aligned} ds^2 &\equiv -c^2 dt^2 + dx^2 + dy^2 + dz^2 \\ &= \eta_{\mu\nu} dx^\mu dx^\nu \end{aligned} \quad \dots \text{pseudo-Riemannian structure}$$

where $x^\mu = (x^0, x^i) = (ct, x, y, z)$ and where

$$\eta_{\mu\nu} = \begin{pmatrix} -1 & & & \\ & +1 & & \\ & & +1 & \\ & & & +1 \end{pmatrix} \equiv \text{diag}(-1, +1, +1, +1) \equiv -1 \oplus \mathbb{1}_3.$$

For photons, $ds^2 = 0$. Using the two postulates, one may show that the interval ds^2 is invariant with respect to coordinates based in any inertial frame (Schutz, 2009, §1.6).

$$ds^2 = \eta_{\mu\nu} dx^\mu dx^\nu = \eta_{\bar{\mu}\bar{\nu}} dx^{\bar{\mu}} dx^{\bar{\nu}}. \quad (\text{I.2})$$

Note

- The symbol ds^2 for the interval is purely notational convention, since we may have $ds^2 < 0$ in some cases.
- The first postulate alone implies either S.R. or its $c \rightarrow \infty$ limit (a.k.a. Galilean relativity).
- Formally, $dx^\mu dx^\nu$ is shorthand for the *tensor product* $dx^\mu \otimes dx^\nu$, and you will see that $ds^2 = \underline{\eta}$ is really the *metric tensor*...

I.1 Lorentz Transformations

In an inertial frame K , the equations of motion of a free particle are linear;

$$\frac{d^2 x^\mu}{d\lambda^2} = 0 \quad \Longleftrightarrow \quad x^\mu = x_0^\mu + u^\mu \lambda,$$

where x_0^μ, u^μ are constant and λ is a parameter. Similarly, in any other inertial frame \bar{K} ,

$$\frac{d^2 x^{\bar{\mu}}}{d\lambda^2} = 0 \quad \Longleftrightarrow \quad x^{\bar{\mu}} = x_0^{\bar{\mu}} + u^{\bar{\mu}} \lambda.$$

Now,

$$\begin{aligned} \frac{dx^{\bar{\mu}}}{d\lambda} &= \frac{\partial x^{\bar{\mu}}}{\partial x^\nu} \frac{dx^\nu}{d\lambda} = \frac{\partial x^{\bar{\mu}}}{\partial x^\nu} u^\nu && \dots \text{chain rule} \\ \Rightarrow 0 &= \frac{d^2 x^{\bar{\mu}}}{d\lambda^2} = \frac{\partial}{\partial x^\alpha} \left(\frac{\partial x^{\bar{\mu}}}{\partial x^\nu} u^\nu \right) u^\alpha = \frac{\partial^2 x^{\bar{\mu}}}{\partial x^\alpha \partial x^\nu} u^\nu u^\alpha && \because u^\nu \text{ constant} \\ \Rightarrow 0 &= \frac{\partial^2 x^{\bar{\mu}}}{\partial x^\alpha \partial x^\nu}. && \dots \text{since true } \forall u^\alpha \end{aligned}$$

So the required transformation between inertial frames is *linear*;

$$\boxed{x^{\bar{\mu}} = L^{\bar{\mu}}{}_\nu x^\nu + a^{\bar{\mu}}}, \quad (\text{I.3})$$

where $a^{\bar{\mu}}$ and $L^{\bar{\mu}}{}_\nu \equiv \frac{\partial x^{\bar{\mu}}}{\partial x^\nu}$ are constants. Differentiate (I.3) and substitute into (I.2) to give $\eta_{\mu\nu} dx^\mu dx^\nu = \eta_{\bar{\mu}\bar{\nu}} L^{\bar{\mu}}{}_\mu L^{\bar{\nu}}{}_\nu dx^\mu dx^\nu$. Since this is true $\forall dx^\mu$,

$$\boxed{\eta_{\bar{\mu}\bar{\nu}} L^{\bar{\mu}}{}_\mu L^{\bar{\nu}}{}_\nu = \eta_{\mu\nu}}. \quad (\text{I.4})$$

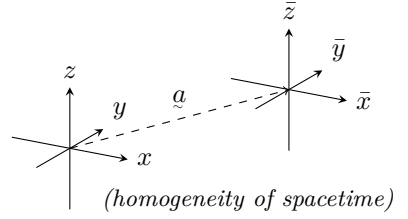
In matrix form, (I.3) and (I.4) are

$$\boxed{\bar{x} = Lx + \underline{a}}, \quad (\text{I.3a})$$

$$L^\top \eta L = \eta. \quad (\text{I.4a})$$

Transformations $x \mapsto \bar{x}$ defined by (I.3), (I.4) are the *inhomogeneous* Lorentz transformations, or **Poincaré transformations**, and form the Poincaré group $\text{IO}(1,3)$ (pronounced Inhomogeneous Orthogonal group).

“Inhomogeneous” refers to the inclusion of spacetime translations $x^\nu \mapsto x^{\bar{\mu}} = \delta^{\bar{\mu}}_\nu x^\nu + a^{\bar{\mu}}$, which form a subgroup T^4 of the Poincaré group. If we set $a^{\bar{\mu}} = 0$ in (I.3), we are left with *homogeneous* transformations, called simply the **Lorentz transformations**.



In G.R., our task is to generalise these ideas to general coordinate frames for which $L^{\bar{\mu}}_\nu$ are not necessarily constant.

I.1.1 Examples of Lorentz Transformations

A transformation belonging to the (homogeneous) Lorentz group $O(1, 3)$ can be represented as a matrix acting on coordinates x^μ when they are *viewed as vectors*.

$$x^\mu \cong \tilde{x} = \begin{pmatrix} x^0 \\ x^1 \\ x^2 \\ x^3 \end{pmatrix} = \begin{pmatrix} t \\ x \\ y \\ z \end{pmatrix}$$

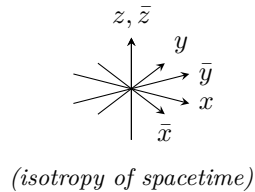
Note

Spacetime events (i.e., *points* in spacetime) are not themselves vectors—neither addition nor scalar multiplication of events makes physical sense (i.e., spacetime itself is not a *vector space*). However, in S.R. we may represent events by their associated displacement vector relative to a chosen orthogonal inertial frame.

The Lorentz group $O(1, 3)$ consists of (combinations of) the following:

- **Rotations**, e.g., about the z -axis (in the xy -plane) by an angle θ ;

$$L_R(\theta) = \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & \cos \theta & -\sin \theta & 0 \\ 0 & \sin \theta & \cos \theta & 0 \\ 0 & 0 & 0 & 1 \end{pmatrix}.$$



- **Boosts**, e.g., by a velocity \vec{v} in the x -direction;

$$L_B(\alpha) = \begin{pmatrix} \cosh \alpha & -\sinh \alpha & 0 & 0 \\ -\sinh \alpha & \cosh \alpha & 0 & 0 \\ 0 & 0 & 1 & 0 \\ 0 & 0 & 0 & 1 \end{pmatrix} = \begin{pmatrix} \cos i\alpha & i \sin i\alpha & 0 & 0 \\ i \sin i\alpha & \cos i\alpha & 0 & 0 \\ 0 & 0 & 1 & 0 \\ 0 & 0 & 0 & 1 \end{pmatrix},$$

where $\alpha = \tanh^{-1} \frac{v}{c}$ is the *rapidity parameter*. In terms of the velocity $\beta \equiv \frac{v}{c}$, one has $\cosh \alpha = \frac{1}{\sqrt{1-\beta^2}} \equiv \gamma$ and $\sinh \alpha = \frac{\beta}{\sqrt{1-\beta^2}} \equiv \beta\gamma$.

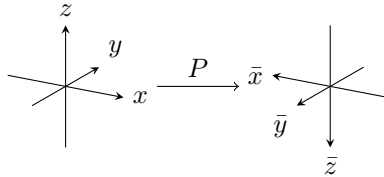
— Note —

Boosts appear similar to rotations, but differ as a consequence of the indefiniteness of the metric. Formally, an x -boost of rapidity α is equivalent to a rotation by an ‘imaginary angle’ $\alpha' = i\alpha$ through the τx -plane, where $\tau = it$ is ‘imaginary time’... though this is not a good picture physically!

An important difference between rotations and boosts is that, where $0 \leq \theta < 2\pi$ for rotations, we have $-\infty < \alpha < \infty$ for boosts, i.e., rotations form a *compact*² subgroup of the Lorentz (or Poincaré) group, whereas boosts are *non-compact*—and in fact do not form a subgroup (because, in general, the composition of two boosts forms a combination of a rotation and a boost).

Both rotations and boosts depend on continuous parameters (θ or α). However, the Lorentz group $O(1, 3)$ also contains *discrete* transformations...

- **Parity inversion**; $P = \text{diag}(1, -1_3) \equiv \begin{pmatrix} 1 & & & \\ & -1 & & \\ & & -1 & \\ & & & -1 \end{pmatrix}$.



Notice that P is *not* equivalent to a rotation; it transforms a right-handed frame into a left and vice versa, since $\det P = -1$.

²A *compact* set is one for which any infinite sequence of elements contains a convergent subsequence. E.g., $[0, 2\pi)$ is compact, but \mathbb{R} is not (consider the sequence $\{1, 2, 3, \dots\} \subset \mathbb{R}$).

Note

In even spatial dimensions, the transformation $R = \text{diag}(1, -\mathbf{1}_{2n})$ is *not* a parity transformation; it is a rotation by π and $\det R = 1$. In these cases, inversions $x_i \mapsto -x_i$ of a *single* spatial coordinate are parity transformations.

- **Time reversal;** $T = \text{diag}(-1, \mathbf{1}_3)$.

The Lorentz matrix condition (I.4) implies that $(\det L)^2 = 1 \iff \det L = \pm 1$ for any Lorentz transformation $L \in \text{O}(1, 3)$. Those with $\det L = +1$ and those with $\det L = -1$ form two disconnected pieces of $\text{O}(1, 3)$, but only the first piece contains the identity transformation δ^μ_ν .

Inspecting the $\bar{\mu}\nu = \bar{0}0$ component of (I.4) gives

$$\left(L^{\bar{0}}_0\right)^2 - \sum_{\bar{k}=1}^3 \left(L^{\bar{k}}_0\right)^2 = 1 \implies \left(L^{\bar{0}}_0\right)^2 \geq 1,$$

which shows that there exists two disconnected classes of Lorentz transformation with $L^{\bar{0}}_0 \geq 1$ and $L^{\bar{0}}_0 \leq -1$. Those with $L^{\bar{0}}_0 \geq 1$ are called **orthochronous**.

I.1.2 The Restricted Lorentz Group

We define the subgroup of restricted Lorentz transformations by adding two conditions to the Lorentz matrix condition (I.4); that they be 1) orthochronous and 2) have determinant unity.

$$\text{SO}^+(1, 3) \equiv \left\{ \Lambda \mid \Lambda^\top \eta \Lambda = \eta, \Lambda^{\bar{0}}_0 \geq 1, \det \Lambda = 1 \right\}$$

We have removed the discrete transformations involving P and T , so that $\text{SO}^+(1, 3)$ is *continuous* and *connected*, unlike $\text{O}(1, 3)$.

Notation

The S in $\text{SO}^+(1, 3)$ refers to the condition $\det \Lambda = 1$, and the $+$ refers to orthochronality. Sometimes $\text{SO}^+(1, 3)$ is simply written as $\text{SO}(1, 3)$.

Groups whose elements may be continuously parametrised are **Lie groups**. The translation group T^4 (parametrised continuously by Δx^μ) and the rotation group $\text{SO}(3)$ (parametrised continuously by three angles) are examples of *connected* Lie groups.

Reintroducing translations to the restricted Lorentz group gives the restricted Poincaré group $\text{ISO}^+(1, 3)$ —also a connected Lie group. Unrestricted groups may be reconstructed by reintroducing the discrete transformations;

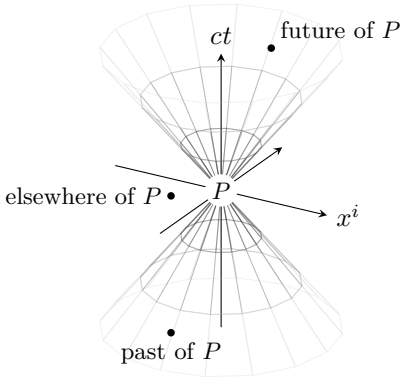
$$\text{O}(1, 3) = \{ \Lambda, \Lambda P, \Lambda T, \Lambda PT \mid \Lambda \in \text{SO}^+(1, 3) \}.$$

I.2 The Scalar Product in Minkowski Space-time

I.3 The Causal Structure of Minkowski Space-time

Because of the indefiniteness of the Minkowski metric (and hence of the Minkowski scalar product), 4-vectors can have positive, zero or negative norm.

A 4-vector \tilde{V} is called $\begin{cases} \text{timelike} \\ \text{null} \\ \text{spacelike} \end{cases}$ if $\tilde{V} \cdot \tilde{V} = V_\mu V^\mu = \eta_{\mu\nu} V^\mu V^\nu \begin{cases} < 0 \\ = 0 \\ > 0 \end{cases}$.



By the second postulate, all *causally connected* events relative to an event P lie in its future or past lightcone, because P cannot causally influence events outside the lightcone without transmitting superluminal data.

A particle's **worldline** is a curve $x^\mu = x^\mu(\lambda)$ whose tangent \tilde{V} given by $V^\mu = \frac{dx^\mu}{d\lambda}$ is everywhere timelike; $V^\mu V_\mu < 0$. This implies that, in an inertial frame, the particles velocity \vec{v} is subluminal $|\vec{v}| < c$ (*exercise*).

One spatial dimension suppressed; lightcone in $(3+1)$ -d spacetime is a continuum of spheres.

exists a Lorentz transformation which sends a timelike unit vector to $V^\mu \mapsto V^{\bar{\mu}} = (1, 0, 0, 0)$ (*exercise*). The **proper time interval** $d\tau$ between two

Consider a freely moving particle whose (timelike) worldline unit tangent vector is \tilde{V} . There always

neighbouring events along the particle's worldline is defined as the interval of time measured in the particle's instantaneous rest frame.

$$c^2 d\tau^2 = -\eta_{\mu\nu} dx^\mu dx^\nu \quad (\text{I.5})$$

Note

Since dx is timelike, $\eta_{\mu\nu} dx^\mu dx^\nu < 0$, so (I.5) contains a minus sign. Proper time is not defined for spacelike intervals, since this would yield an imaginary time. Instead, proper distance is defined by $d\ell^2 = \eta_{\mu\nu} dx^\mu dx^\nu$.

Proper time is the time *experienced* by the particle, i.e., the time that would be measured by a clock moving on the same worldline.

Suppose we have a worldline $x^\mu(\lambda)$. Integrate $d\tau$ via (I.5) along the path to get the total proper time elapsed between events $x^\mu(\lambda_i)$ and $x^\mu(\lambda_f)$.

$$\Delta\tau = \frac{1}{c} \int_{\lambda_i}^{\lambda_f} d\lambda \sqrt{-\eta_{\mu\nu} \frac{dx^\mu}{d\lambda} \frac{dx^\nu}{d\lambda}} \quad (\text{I.6})$$

Note

The proper time elapsed depends on the path taken between the two events. (See the Twin Paradox.)

Particle worldlines can be characterised by an action principle: variation of (I.6) with respect to the trajectories, $\delta x^\mu(\lambda)$, yields the equations of motion $\frac{\partial^2 x^\mu}{\partial \lambda^2} = 0$ (holding $\delta x^\mu = 0$ fixed at both ends) (*exercise*). In fact, the extremised trajectory minimises the proper time—*freely falling particles take the path of maximum proper time*. □

By (I.5), $c^2 d\tau^2 = -dt^2 + \sum_i (dx^i)^2 = dt^2(-1 + \vec{v} \cdot \vec{v})$ where $v^i \equiv \frac{dx^i}{dt}$ so

$$d\tau = \sqrt{1 - \frac{v^2}{c^2}} dt \equiv \gamma dt \quad (\text{I.7})$$

is the proper time of a frame moving at velocity \vec{v} relative to a frame with coordinates (ct, x^i) .

I.3.1 4-velocity

Since proper time $d\tau$ is an *invariant*, we are motivated to define the 4-vector

$$u^\mu \equiv \frac{dx^\mu}{d\tau} = \frac{dt}{d\tau} \left(c, \frac{d\vec{x}}{dt} \right) = \gamma(c, \vec{v}),$$

where γ is defined as in (I.7). In the particle's instantaneous rest frame, $\|\vec{v}\|^2 = u^i u_i = 0$, so that $u^\mu u_\mu = u^0 u_0 = -c^2$. However, since this is an invariant (it is a scalar, which does not transform under Poincaré transformations), we have

$$\boxed{\|\underline{u}\|^2 \equiv \underline{u} \cdot \underline{u} \equiv u^\mu u_\mu = -c^2}$$

in *all frames*.

I.3.2 4-momentum

Relativistic 4-momentum is defined by

$$\underline{p} \equiv m \underline{u} = \left(\frac{E}{c}, \vec{p} \right), \quad E = \gamma m c^2, \quad \vec{p} = \gamma m \vec{v}. \quad (\text{I.8})$$

This form of 4-momentum may be deduced from an action using the Hamiltonian approach. We want to extremise (I.2), but first multiply by $m c^2$ to get the units right. The action is

$$S = -m c^2 \int_{x_i}^{x_f} d\tau = \int_{t_i}^{t_f} dt L$$

where the Lagrangian is

$$L = L(\vec{x}, \dot{\vec{x}}; t) = -m c^2 \sqrt{1 - \frac{v^2}{c^2}} = -m c^2 \sqrt{1 - \frac{\vec{x} \cdot \dot{\vec{x}}}{c^2}}.$$

Then, the Hamiltonian (total energy) and conjugate momentum are

$$\begin{aligned} E &= \vec{p} \cdot \dot{\vec{x}} - L = \gamma m v^2 + \frac{m c^2}{\gamma} = \gamma m c^2, \\ \vec{p} &= \frac{\partial L}{\partial \dot{\vec{x}}} = m c^2 \frac{1}{2} \frac{2 \dot{\vec{x}}}{\sqrt{1 - v^2/c^2}} = \gamma m \vec{v}, \end{aligned}$$

matching the definition (I.8).

In all rest frames,

$$\boxed{\|\underline{p}\|^2 = \underline{p} \cdot \underline{p} = p^\mu p_\mu = -m^2 c^2}$$

as can be shown by evaluating in the particle's instantaneous rest frame and using scalar invariance.

For photons with $m = 0$, the 4-momentum p is a null vector. Photons have no instantaneous rest frame—but we can always find a frame in which $p = \frac{E}{c}(1, 0, 0, 1)$.

I.3.3 Angular momentum

The covariant orbital angular momentum of a particle about an event $x_0 = \{x_0^\mu\}$ is defined as

$$j^{\mu\nu} = (x^\mu - x_0^\mu)p^\nu - (x^\nu - x_0^\nu)p^\mu,$$

where x^μ and p^μ are the particle's position and momentum. This is an antisymmetric tensor, $j^{\mu\nu} = -j^{\nu\mu} = j^{[\mu\nu]}$, with 6 independent components: $j^{0k}, j^{k\ell}, k < \ell$.

— *Example* —

About the origin $x_0^\mu = (0, 0, 0, 0)$, we have $j^{ij} = x^i p^j - x^j p^i$ which is equivalent to $j^{ij} = \epsilon^{ijk} \ell_k$ where $\vec{\ell} = \vec{x} \times \vec{p}$. j^{0k} represents the motion of the centre of mass: $j^{0k} = x^0 p^k - x^k p^0 = \gamma mc(v^k t - x^k)$.

I.4 Gauß's Theorem in Minkowski Spacetime

II The Equivalence Principle

References

- N. Doughty. *Lagrangian interaction: an introduction to relativistic symmetry in electrodynamics and gravitation*. CRC Press, 2018.
- B. Schutz. *A first course in general relativity*. Cambridge university press, 2009.