11/2 Lecture outline

⋆ **Reading for this week's lecture: Coleman lecture notes pages 109-139.**

• Recall examples of $2 \to 2$ processes, e.g. $N+N \to N+N$, to $\mathcal{O}(g^2)$

$$N+N \to N+N: \qquad i\mathcal{A} = (-ig)^2 \left(\frac{i}{t-\mu^2+i\epsilon} + \frac{1}{u-\mu^2+i\epsilon} \right)$$

$$N+\bar{N} \to N+\bar{N}: \qquad i\mathcal{A} = (-ig)^2 \left(\frac{i}{t-\mu^2+i\epsilon} + \frac{1}{s-\mu^2+i\epsilon} \right)$$

where $s=(p_1+p_2)^2$, $t=(p_1-p_1')^2$, $u=(p_1-p_2')^2$, with $s+t+u=4m^2$ (more generally, $s+t+u=\sum_{i=1}^4 m_i^2$). In CM, $s=4E^2$, $t=-2p^2(1-\cos\theta)$, and $u=-2p^2(1+\cos\theta)$.

• Crossing symmetry, CPT. Write $1+2 \to \bar{3}+\bar{4}$. Take all momenta incoming, $\mathcal{A}(p_1,p_2,p_3,p_4)$, with $p_1+p_2+p_3+p_4=0$ and use $s=(p_1+p_2)^2$, $t=(p_1+p_3)^2$ and $u=(p_1+p_4)^2$. Note $s+t+u=\sum_{n=1}^4 m_n^2$. The process $1+2\to\bar{3}+\bar{4}$ is kinematically allowed for $s>4m^2$, $t<0$, $u<0$. If instead $u>4m^2$, it's the process $1+3\to\bar{2}+\bar{4}$.

• Scattering by ϕ exchange leads to an attractive Yukawa potential. This was Yukawa's original goal, to explain the attraction between nucleons. Indeed, the t-channel term in e.g. the above $N+N$ scattering amplitude gives, upon using $(p_1-p_1')^2-\mu^2=-(|\vec{p}_1-\vec{p}_1'|^2+\mu^2)$, and the Born approximation[1] in NRQM, $\mathcal{A}_{NR}=-\int d^3\vec{r}e^{-i(\vec{p}'-\vec{p})\cdot\vec{r}}V(\vec{r})$, the attractive Yukawa potential

$$V(r)=\int \frac{d^3q}{(2\pi)^3}\frac{-(g/2m)^2 e^{i\vec{q}\cdot\vec{r}}}{|\vec{q}|^2+\mu^2} = -\frac{(g/2m)^2}{4\pi r}e^{-\mu r}.$$

(The $1/(2m)^2$ is because we normalized the relativistic states with the extra factor of $2E\approx 2m$ as compared with standard nonrelativistic normalization[2]. This gives Yukawa's explanation of the attraction between nucleons, from meson exchange. The u-channel term is an exchange potential interaction, which exchanges the positions of the two identical particles in addition to giving a potential. For angular momentum ℓ in a partial-wave expansion, the exchange term differs from the direct one by a factor of $(-1)^\ell$.

[1] Max Born, in QM, or Lord Rayleigh classically: $\frac{d\sigma}{d\Omega}\sim |U(\vec{q})|^2$

[2] This is clear on dimensional grounds, since $[g]\sim m$. More generally, write $a(p)=\sqrt{2E}\widehat{a}(p)$ and $\mathcal{A}=\prod_i\sqrt{2E_i}\prod_f\sqrt{2E_f}\widehat{\mathcal{A}}$.

• We saw above that the t channel term above is associated with the Yukawa potential. The u channel term is similar. Now consider the s channel, in e.g. the $N + \bar{N}$ scattering amplitude. Using the CM relations $\vec{p}_1 = -\vec{p}_2 \equiv \vec{p}$ and $E_1 = E_2 = \sqrt{p^2 + m^2}$ gives

$$\mathcal{A} \sim \frac{1}{4m^2 + 4p^2 - \mu^2 + i\epsilon},$$

so for $\mu < 2m$ the denominator is always positive, and the amplitude decreases with increasing p^2. For $\mu > 2m$ there is a pole at $(p_1 + p_2)^2 = \mu^2$, where the intermediate meson goes on shell. This leads to a peak (not a pole, of course; because the intermediate particle is unstable anyway, the denominator gets an imaginary contribution from higher order contributions), a *resonance*, in the cross section. E.g. Z_0 pole in $e^+e^- \to \mu^+\mu^-$, but not in $e^+e^- \to \gamma\gamma$.

• Solve $\mathcal{L} = \frac{1}{2}\partial\phi^2 - \frac{1}{2}m^2\phi^2 - J(x)\phi$. Using Dyson + Wick's theorem, $U(\infty, -\infty) =:$ $e^{O_1 + \frac{1}{2}O_2}$:, where $O_1 = -i\int d^4x J(x)\phi(x)$ and $O_2 = (-i)^2 \int d^4x_1 d^4x_2 D_F(x_1 - x_2) J(x_1) J(x_2)$. So $O_2 = \alpha + i\beta$ is a number, whereas O_1 is an operator. Will lead to probability P_n for creating out of the vacuum a state with n mesons given by $P_n = e^{-|\alpha|}|\alpha|^n/n!$, the Poisson distribution. You'll work out the details in the HW assignment.

• Compute probabilities by squaring the S-maxtrix amplitudes, but have to be careful with the delta functions, since squaring the delta functions would give nonsense.

Warmup: consider quantum mechanics, with $U(t) = Te^{-i\int^t H(t)dt}$,

$$\langle f|U(t)|i\rangle \approx -i\langle f|H_{int}|i\rangle \int_0^t dt e^{i\omega t},$$

where $\omega = E_f - E_i$. If we take $t \to \infty$ first, we get $\delta(\omega)$ and squaring would give nonsense. That's because we're asking the wrong question if we ask about probability for a transition over all time – instead, we should ask about the rate. So keep t finite for now. Squaring gives $P(t) = 2|\langle f|H_{int}|i\rangle|^2(1 - \cos\omega t)/\omega^2$. For $t \to \infty$, multiply by $dE_f\rho(E_f)$ and replace $(1 - \cos\omega t)/\omega^2 = 4\sin^2(\frac{1}{2}\omega t)/\omega^2 \to \pi t\delta(\omega)$ (using $\int_{-\infty}^{\infty} dx x^{-2}\sin^2 x = \pi$ (hint: $\sin^2 x/x^2 = (2 - e^{i2x} - e^{-i2x})/4x^2$ and close the contour in the correct directions)) to get

$$\dot{P}_{i\to f} = 2\pi|\langle f|H_{int}|i\rangle|^2\rho(E).$$

This is "Fermi's Golden Rule" – it was actually derived by Dirac, but Fermi used it a lot and called it the golden rule. Another aside: Fermi and Dirac independently discovered that spin 1/2 objects must anticommute, and Dirac generously named such objects "Fermions".

Naively taking $t \to \infty$ initially would have given amplitude $\sim \delta(\omega)$, and squaring that would give $\delta(\omega)^2$, which needs to be replaced with $\delta(\omega) 2\pi T$, and then divide by T to get the rate. Similarly in field theory, $\delta(p)^2$ should be replaced with probability $\sim \delta(p)$ times phase space volume factors.

• Phase space factors. Put the system in a box of volume V. The momenta are quantized and, as usual, there are $V d^3\vec{k}/(2\pi)^3$ states with $\vec{k}$ in the range $d^3\vec{k}$. Continue next time..