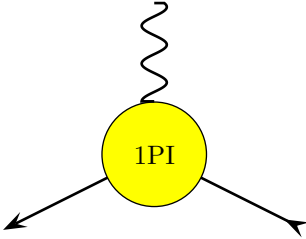
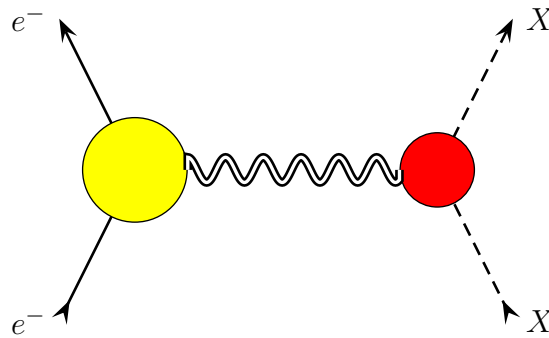


QED Vertex Correction

Consider the dressed electron-electron-photon vertex in QED,

$$ie\Gamma^\mu(p', p) = \text{1PI} \quad (1)$$


We are interested in this vertex in the context of elastic Coulomb scattering,



$$\quad (2)$$

so we take the incoming and the outgoing electrons to be on-shell, $p^2 = p'^2 = m^2$, but the photon is off-shell, $q^2 \neq 0$. Moreover, we put the vertex in the context of the complete electron line — including the external line factors, thus $\bar{u}(p') \times ie\Gamma^\mu \times u(p)$. As we saw in the [my notes on the form factors](#), this simplifies the Lorentz and Dirac structure of the vertex and allows us to write it as

$$\Gamma^\mu(p', p) = F_{\text{el}}(q^2) \times \frac{(p' + p)^\mu}{2m} + F_{\text{mag}}(q^2) \times \frac{i\sigma^{\mu\nu}q_\nu}{2m} = F_1(q^2) \times \gamma^\mu + F_2(q^2) \times \frac{i\sigma^{\mu\nu}q_\nu}{2m}. \quad (3)$$

At the tree level, the electron is a point-like spin = $\frac{1}{2}$ particle obeying Dirac equations, hence $F_1(q^2) \equiv 1$ and $F_2(q^2) \equiv 0$. But the quantum corrections in QED mix the elementary electron state with the composite states like $|e^- \gamma\rangle$, $|e^- e^- e^+\rangle$, etc., and this leads to the non-trivial q^2 -dependent form-factors.

In these notes we shall calculate the $F_1(q^2)$ and the $F_2(q^2)$ form-factors to the one-loop order in QED. But before we draw and evaluate specific loop diagrams, note that in the counterterm

perturbation theory

Diagram (4) shows a yellow circle labeled 'net' with a wavy line entering from the top and two straight lines exiting downwards. This is equal to the sum of three diagrams: a tree-level vertex (black dot), a yellow circle labeled 'loops' (identical to the 'net' diagram), and a pink circle (representing a counterterm) with a wavy line entering from the top and two straight lines exiting downwards.

thus

$$ie\Gamma_{\text{net}}^\mu(p', p) = (ie\gamma^\mu)_{\text{tree}} + ie\Gamma_{\text{loops}}^\mu(p', p) + ie\delta_1 \times \gamma^\mu, \quad (5)$$

where the counterterm coefficient δ_1 is set such that

$$\Gamma_{\text{net}}^\mu(\text{on-shell } p' = p) = \gamma^\mu, \text{ exactly.} \quad (6)$$

For on-shell incoming and outgoing electrons but off-shell photon, eq. (5) translates to the counterterm language as

$$\begin{aligned} F_1^{\text{net}}(q^2) &= 1^{\text{tree}} + F_1^{\text{loops}}(q^2) + \delta_1, \\ F_2^{\text{net}}(q^2) &= 0^{\text{tree}} + F_2^{\text{loops}}(q^2), \end{aligned} \quad (7)$$

while eq. (6) becomes

$$\delta_1 = -F_1^{\text{loops}}(q^2 = 0) \implies F_1^{\text{net}}(q^2 = 0) = 1. \quad (8)$$

One-Loop Calculation: Working Through the Algebra

Now let's turn to the actual calculation of the dressed QED vertex Γ_{net}^μ to the order $O(\alpha)$. Fortunately for us, at this order there is only one 1-loop diagram dressing up the vertex, namely

Diagram (9) shows a yellow circle labeled '1 loop' with a wavy line entering from the top and two straight lines exiting downwards. This is equal to a diagram with a wavy line entering from the top, a black dot at the top vertex, and two straight lines exiting downwards. A wavy line loop connects the two lower vertices, with black dots at the vertices where the loop meets the external lines.

In the Feynman gauge for the internal photon's propagator, this diagram evaluates to

$$\begin{aligned}
ie\Gamma_{1\text{loop}}^\mu(p', p) &= \int_{\text{reg}} \frac{d^4k}{(2\pi)^4} \frac{-ig^{\nu\lambda}}{k^2 + i0} \times ie\gamma_\nu \times \frac{i}{\not{p}' + \not{k} - m + i0} \times ie\gamma^\mu \times \frac{i}{\not{p} + \not{k} - m + i0} \times ie\gamma_\lambda \\
&= e^3 \int_{\text{reg}} \frac{d^4k}{(2\pi)^4} \frac{1}{k^2 + i0} \times \gamma^\nu \times \frac{\not{p}' + \not{k} + m}{(p' + k)^2 - m^2 + i0} \times \gamma^\mu \times \frac{\not{p} + \not{k} + m}{(p + k)^2 - m^2 + i0} \times \gamma_\nu \\
&= e^3 \int_{\text{reg}} \frac{d^4k}{(2\pi)^4} \frac{\mathcal{N}^\mu}{\mathcal{D}}
\end{aligned} \tag{10}$$

where

$$\mathcal{N}^\mu = \gamma^\nu (\not{k} + \not{p}' + m) \gamma^\mu (\not{k} + \not{p} + m) \gamma_\nu \tag{11}$$

and

$$\mathcal{D} = [k^2 + i0] \times [(p + k)^2 - m^2 + i0] \times [(p' + k)^2 - m^2 + i0]. \tag{12}$$

The purpose of this section of the notes is to simplify these numerator and denominator. Using the Feynman parameter trick, we may combine the 3 denominator factors as

$$\frac{1}{\mathcal{D}} = \iiint_0^1 dx dy dz \delta(x + y + z - 1) \frac{2}{\left[x((p + k)^2 - m^2) + y((p' + k)^2 - m^2) + z(k^2) + i0 \right]^3}. \tag{13}$$

Inside the big square brackets here we have

$$\begin{aligned}
[\dots] &= x \times ((p + k)^2 - m^2) + y \times ((p' + k)^2 - m^2) + z \times k^2 \\
&= k^2 \times (x + y + z = 1) + 2k_\mu (xp + yp')^\mu + x(p^2 - m^2) + y(p'^2 - m^2) \\
&= (k + xp + yp')^2 - \Delta
\end{aligned} \tag{14}$$

where

$$\Delta = (xp + yp')^2 - xp^2 - yp'^2 + (x + y)m^2. \tag{15}$$

In this formula

$$\begin{aligned}
(xp + yp')^2 &= x^2p^2 + y^2p'^2 + xy \times (2p \cdot p' = p^2 + p'^2 - (p' - p)^2) \\
&= x(x + y)p^2 + y(y + x)p'^2 - xyq^2,
\end{aligned} \tag{16}$$

hence

$$\begin{aligned}\Delta &= x(x+y-1)p^2 + y(x+y-1)p'^2 - xyq^2 + (x+y)m^2 \\ &= -xzp^2 - yzp'^2 - xyq^2 + (1-z)m^2.\end{aligned}\tag{17}$$

Moreover, for the on-shell electron momenta, $p^2 = p'^2 = m^2$, we may further simplify

$$(1-z) \times m^2 - xz \times p^2 - yz \times p'^2 = m^2 \times \left((1-z) - (x+y=1-z)z = (1-z)^2 \right),\tag{18}$$

hence

$$\Delta = (1-z)^2 \times m^2 - xy \times q^2.\tag{19}$$

Let us also define the shifted loop momentum

$$\ell = k + xp + yp',\tag{20}$$

then we can rewrite the denominator as

$$\frac{1}{\mathcal{D}} = \iiint_0^1 dx dy dz \delta(x+y+z-1) \frac{2}{[\ell^2 - \Delta + i0]^3}.\tag{21}$$

As usual, we plug this denominator into the loop integral (10), then change the order of integration — \int over the loop momentum before \int over the Feynman parameters, — and then shift the momentum integration variable from k to ℓ , thus

$$\Gamma_{1\text{loop}}^\mu(p', p) = -2ie^2 \iiint_0^1 dx dy dz \delta(x+y+z-1) \int_{\text{reg}} \frac{d^4\ell}{(2\pi)^4} \frac{\mathcal{N}^\mu}{[\ell^2 - \Delta + i0]^3}.\tag{22}$$

But to make full use of the momentum shift, we need to re-express the numerator \mathcal{N}^μ in terms of the shifted momentum ℓ . It would also help to simplify the numerator (11) *in the context of this monstrous integral*.

The first step towards simplifying the \mathcal{N}^μ is obvious: Let us get rid of the γ^ν and γ_ν factors using the γ matrix algebra, *eg.*, $\gamma^\nu \not{a} \gamma_\nu = -2 \not{a}$, *etc.*. However, in order to allow for the dimensional regularization, we need to re-work the algebra for an arbitrary spacetime dimension D where $\gamma^\nu \gamma_\nu = D \neq 4$. So let X be some product of n Dirac matrices, then

$$\gamma^\nu X \gamma_\nu = [\gamma^\nu, X] \gamma_\nu + (-1)^n \times X \times (\gamma^\nu \gamma_\nu = D), \quad (23)$$

where the (anti)commutator term is D independent and only the second term is affected by $D \neq 4$. Hence,

$$\left(\gamma^\nu X \gamma_\nu \right)_{D \neq 4} = \left(\gamma^\nu X \gamma_\nu \right)_{4d} + (-1)^n (D - 4) \times X; \quad (24)$$

in particular, for $n = 1, 2, 3$ we have

$$\begin{aligned} \gamma^\nu \not{a} \gamma_\nu &= -2 \not{a} + (4 - D) \times \not{a}, \\ \gamma^\nu \not{a} \not{b} \gamma_\nu &= 4(ab) - (4 - D) \times \not{a} \not{b}, \\ \gamma^\nu \not{a} \not{b} \not{c} \gamma_\nu &= -2 \not{c} \not{b} \not{a} + (4 - D) \times \not{a} \not{b} \not{c}. \end{aligned} \quad (25)^*$$

Thus, applying these formulae to the numerator (11) we arrive at

$$\begin{aligned} \mathcal{N}^\mu &\stackrel{\text{def}}{=} \gamma^\nu (\not{k} + \not{p}' + m) \gamma^\mu (\not{k} + \not{p} + m) \gamma_\nu \\ &= -2m^2 \gamma^\mu + 4m(p' + p + 2k)^\mu - 2(\not{p} + \not{k}) \gamma^\mu (\not{p}' + \not{k}) \\ &\quad + (4 - D)(\not{p}' + \not{k} - m) \gamma^\mu (\not{p} + \not{k} - m). \end{aligned} \quad (26)$$

The second step is to re-express this numerator in terms of the loop momentum ℓ rather than k using eq. (20). Expanding the result in powers of ℓ , we get quadratic, linear and ℓ -independent terms, but the linear terms do not contribute to the $\int d^D \ell$ integral because they are odd with respect to $\ell \rightarrow -\ell$ while everything else in that integral is even. Consequently, *in the context of*

eq. (22) we may neglect the linear terms, thus

$$\begin{aligned}
\mathcal{N}^\mu &= -2m^2\gamma^\mu + 4m(p' + p + 2\ell - 2xp - 2yp')^\mu \\
&\quad - 2(\not{p} + \not{\ell} - x\not{p} - y\not{p}')\gamma^\mu(\not{p}' + \not{\ell} - x\not{p} - y\not{p}') \\
&\quad + (4 - D)(\not{p}' + \not{\ell} - x\not{p} - y\not{p}' - m)\gamma^\mu(\not{p} + \not{\ell} - x\not{p} - y\not{p}' - m) \\
&\quad \langle\langle \text{skipping terms linear in } \ell \rangle\rangle \\
&\cong -2m^2\gamma^\mu + 4m(p + p' - 2xp - 2yp')^\mu \\
&\quad - 2\not{\ell}\gamma^\mu\not{\ell} - 2(\not{p} - x\not{p} - y\not{p}')\gamma^\mu(\not{p}' - x\not{p} - y\not{p}') \\
&\quad + (4 - D)\not{\ell}\gamma^\mu\not{\ell} + (4 - D)(\not{p}' - y\not{p}' - x\not{p} - m)\gamma^\mu(\not{p} - x\not{p} - y\not{p}' - m).
\end{aligned} \tag{27}$$

Next, we make use of $p' - p = q$ and $1 - x - y = z$ to rewrite

$$\begin{aligned}
2xp + 2yp' &= (x + y) \times (p + p') + (x - y) \times (p - p'), \\
p + p' - 2xp - 2yp' &= z \times (p' + p) + (x - y) \times q, \\
p - xp - yp' &= z \times p - y \times q \\
&= z \times p' - (1 - x) \times q, \\
p' - xp - yp' &= z \times p' + x \times q \\
&= z \times p + (1 - y) \times q,
\end{aligned} \tag{28}$$

and consequently

$$\begin{aligned}
\mathcal{N}^\mu &\cong -2m^2\gamma^\mu + 4mz(p' + p)^\mu + 4m(x - y)q^\mu \\
&\quad + (-2 + 4 - D) \times \not{\ell}\gamma^\mu\not{\ell} \\
&\quad - 2(z\not{p}' - (1 - x)\not{q})\gamma^\mu(z\not{p} + (1 - y)\not{q}) \\
&\quad + (4 - D)(z\not{p}' + x\not{q} - m)\gamma^\mu(z\not{p} - y\not{q} - m).
\end{aligned} \tag{29}$$

The third step is to make use of the external fermions being on-shell. This means more than just $p^2 = p'^2 = m^2$: We also sandwich the vertex $ie\Gamma^\mu$ between the Dirac spinors $\bar{u}(p')$ on the left and $u(p)$ on the right. The two spinors satisfy the appropriate Dirac equations $\not{p}u(p) = mu(p)$ and $\bar{u}(p')\not{p}' = \bar{u}(p')m$, so in the context of $\bar{u}(p')\Gamma^\mu u(p)$,

$$A \times \not{p} \cong A \times m \quad \text{and} \quad \not{p}' \times B \cong m \times B \tag{30}$$

for any terms in Γ^μ that look like $A \times \not{p}$ or $\not{p}' \times B$ for some A or B . Consequently, the terms on

the last two lines of eq. (29) are equivalent to

$$\begin{aligned}
(z \not{p}' - (1-x) \not{q}) \gamma^\mu (z \not{p} + (1-y) \not{q}) &\cong (zm - (1-x) \not{q}) \gamma^\mu (zm + (1-y) \not{q}) \\
&= z^2 m^2 \times \gamma^\mu - (1-x)(1-y) \times \not{q} \gamma^\mu \not{q} \\
&\quad + z(x-y)m \times \left(\frac{1}{2} \{ \gamma^\mu, \not{q} \} = q^\mu \right) \\
&\quad + z(2-x-y)m \times \left(\frac{1}{2} [\gamma^\mu, \not{q}] = -i\sigma^{\mu\nu} q_\nu \right) \quad (31)
\end{aligned}$$

where

$$z(2-x-y) = z(1+z), \quad (32)$$

$$(1-x)(1-y) = 1-x-y+xy = z+xy, \quad (33)$$

and $(z \not{p}' + x \not{q} - m) \gamma^\mu (z \not{p} - y \not{q} - m) \cong ((z-1)m + x \not{q}) \gamma^\mu ((z-1)m - y \not{q})$

$$\begin{aligned}
&= (1-z)^2 m^2 \times \gamma^\mu - xy \times \not{q} \gamma^\mu \not{q} \\
&\quad - (1-z)(x-y)m \times \left(\frac{1}{2} \{ \gamma^\mu, \not{q} \} = q^\mu \right) \\
&\quad + (1-z)(x+y)m \times \left(\frac{1}{2} [\gamma^\mu, \not{q}] = -i\sigma^{\mu\nu} q_\nu \right) \quad (34)
\end{aligned}$$

where

$$(1-z)(x+y) = (1-z)^2. \quad (35)$$

Let's plug these expressions back into eq. (29) and collect similar terms together, thus

$$\begin{aligned}
\mathcal{N}^\mu &\cong -(D-2) \not{q} \gamma^\mu \not{q} + 4mz(p' + p)^\mu \\
&\quad + m^2 \gamma^\mu \times \left(-2 - 2z^2 + (4-D)(1-z)^2 \right) \\
&\quad + \not{q} \gamma^\mu \not{q} \times \left(2(z+xy) - (4-D)xy \right) \quad (36) \\
&\quad + mq^\mu \times (x-y) \left(4 - 2z - (4-D)(1-z) \right) \\
&\quad + im\sigma^{\mu\nu} q_\nu \times \left(2z(1+z) - (4-D)(1-z)^2 \right).
\end{aligned}$$

Furthermore, in the context of the Dirac sandwich $\bar{u}(p') \Gamma^\mu u(p)$ we have

$$\not{q} \gamma^\mu \not{q} = 2q^\mu \not{q} - q^2 \gamma^\mu \cong -q^2 \gamma^\mu \quad (37)$$

because $\bar{u}(p') \not{q} u(p) = 0$, and also the Gordon identity

$$(p' + p)^\mu \cong 2m\gamma^\mu - i\sigma^{\mu\nu} q_\nu. \quad (38)$$

Plugging these formulae into eq. (36) and collecting similar terms, we arrive at

$$\begin{aligned}
\mathcal{N}^\mu &\cong -(D-2) \not{\ell} \gamma^\mu \not{\ell} + m^2 \gamma^\mu \times \left(8z - 2(1+z^2) + (4-D)(1-z)^2 \right) \\
&\quad - q^2 \gamma^\mu \times \left(2(z+xy) - (4-D)xy \right) - im\sigma^{\mu\nu} q_\nu \times \left(4z - 2z(1+z) + (4-D)(1-z)^2 \right) \\
&\quad + mq^\mu \times (x-y) \left(4 - 2z - (4-D)(1-z) \right).
\end{aligned} \tag{39}$$

To further simplify this expression, let us go back to the symmetries of the integral (22). The integral over the Feynman parameters, the integral $\int d^D \ell$, and the denominator $[l^2 - \Delta]^3$ are all invariant under the parameter exchange $x \leftrightarrow y$. In eq. (39) for the numerator, the first two lines are invariant under this symmetry, but the last line changes sign. Consequently, only the first two lines contribute to the integral (22) while the third line integrates to zero and may be disregarded, thus

$$\begin{aligned}
\mathcal{N}^\mu &\cong -(D-2) \not{\ell} \gamma^\mu \not{\ell} + m^2 \gamma^\mu \times \left(8z - 2(1+z^2) + (4-D)(1-z)^2 \right) \\
&\quad - q^2 \gamma^\mu \times \left(2(z+xy) - (4-D)xy \right) - im\sigma^{\mu\nu} q_\nu \times \left(4z - 2z(1+z) + (4-D)(1-z)^2 \right) \\
&= -(D-2) \not{\ell} \gamma^\mu \not{\ell} + m^2 \gamma^\mu \times \left(4z - (D-2)(1-z)^2 \right) \\
&\quad - q^2 \gamma^\mu \times \left(2z + (D-2)xy \right) - im\sigma^{\mu\nu} q_\nu \times (1-z) \left(2z + (4-D)(1-z) \right).
\end{aligned} \tag{40}$$

Moreover, thanks to the Lorentz invariance of the $\int d^D \ell$ integral,

$$\ell_\lambda \ell_\nu \cong g_{\lambda\nu} \times \frac{\ell^2}{D}, \tag{41}$$

hence

$$\not{\ell} \gamma^\mu \not{\ell} = \gamma^\lambda \gamma^\mu \gamma^\nu \times \ell_\lambda \ell_\nu \cong \gamma^\lambda \gamma^\mu \gamma^\nu \times g_{\lambda\nu} \frac{\ell^2}{D} = -(D-2) \gamma^\mu \times \frac{\ell^2}{D}, \tag{42}$$

and therefore

$$\begin{aligned}
\mathcal{N}^\mu &\cong + \frac{(D-2)^2}{D} \times \ell^2 \gamma^\mu + m^2 \gamma^\mu \times \left(4z - (D-2)(1-z)^2 \right) \\
&\quad - q^2 \gamma^\mu \times \left(2z + (D-2)xy \right) - im\sigma^{\mu\nu} q_\nu \times (1-z) \left(2z + (4-D)(1-z) \right).
\end{aligned} \tag{43}$$

Finally, note the γ -matrix structure of the 4 terms on the RHS of eq. (43): the first 3 terms are proportional to the γ^μ , so their integrals contribute to the $F_1(q^2)$ form factor, while the fourth

term is proportional to the $-i\sigma^{\mu\nu}q_\nu$ so its integral contributes to the $F_2(q^2)$. Indeed, reorganizing the numerator (43) as

$$\mathcal{N}^\mu \cong \mathcal{N}_1 \times \gamma^\mu - \mathcal{N}_2 \times \frac{i\sigma^{\mu\nu}q_\nu}{2m} \quad (44)$$

$$\begin{aligned} \text{where } \mathcal{N}_1 &= \frac{(D-2)^2}{D} \times \ell^2 + \left(4z - (D-2)(1-z)^2\right) \times m^2 \\ &\quad - \left(2z + (D-2)xy\right) \times q^2 \\ &= \frac{(D-2)^2}{D} \times \ell^2 - (D-2) \times \Delta + 2z \times (2m^2 - q^2) \end{aligned} \quad (45)$$

$$\text{and } \mathcal{N}_2 = 2(1-z) \left(2z + (4-D)(1-z)\right) \times m^2 \quad (46)$$

and plugging this reorganization into the integral (22), we arrive at

$$\Gamma_{1\text{loop}}^\mu = F_1^{1\text{loop}}(q^2) \times \gamma^\mu + F_2^{1\text{loop}}(q^2) \times \frac{i\sigma^{\mu\nu}q_\nu}{2m}, \quad (47)$$

$$F_1^{1\text{loop}}(q^2) = -2ie^2 \iiint_0^1 dx dy dz \delta(x+y+z-1) \int \frac{d^D\ell}{(2\pi)^D} \frac{\mathcal{N}_1}{[\ell^2 - \Delta + i0]^3}, \quad (48)$$

$$F_2^{1\text{loop}}(q^2) = +2ie^2 \iiint_0^1 dx dy dz \delta(x+y+z-1) \int \frac{d^D\ell}{(2\pi)^D} \frac{\mathcal{N}_2}{[\ell^2 - \Delta + i0]^3}. \quad (49)$$

Electron's Gyromagnetic Moment

As explained earlier in class, electron's spin couples to the static magnetic field as

$$\hat{H} \supset \frac{-eg}{2m_e} \mathbf{S} \cdot \mathbf{B} \quad \text{where } g = 2 \left(F_{\text{mag}} = F_1 + F_2 \right) \Big|_{q^2=0}. \quad (50)$$

Moreover, for $q^2 = 0$ the electric form factor $F_1 \equiv F_{el}$ is constrained by the net charge conservation,

$$F_1^{\text{tot}} = 1^{\text{tree}} + F_1^{\text{loops}} + \delta_1 \xrightarrow{q^2 \rightarrow 0} 1. \quad (51)$$

Therefore, the gyromagnetic moment is

$$g = 2 + 2F_2(q^2 = 0) \quad (52)$$

where $F_2 = F_2^{\text{loops}}$ because there are no tree-level or counter-term contributions to the F_2 ,

only to the F_1 . Thus, to calculate the $g - 2$ at the one-loop level, all we need is to evaluate the integral (49) for $q^2 = 0$.

Let's start with the momentum integral

$$\int \frac{d^D \ell}{(2\pi)^D} \frac{\mathcal{N}_2}{[\ell^2 - \Delta + i0]^3} \quad (53)$$

where $\Delta = (1 - z)^2 m^2$ for $q^2 = 0$ and \mathcal{N}_2 is as in eq. (46). Because the numerator here does not depend on the loop momentum ℓ , this integral converges in $D = 4$ dimensions and there is no need for dimensional regularization. All we need is to rotate the momentum into Euclidean space,

$$\begin{aligned} \int \frac{d^4 \ell}{(2\pi)^4} \frac{\mathcal{N}_2}{[\ell^2 - \Delta + i0]^3} &= \mathcal{N}_2 \times \int \frac{i d^4 \ell_E}{(2\pi)^4} \frac{1}{-(\ell_E^2 + \Delta)^3} \\ &= \frac{-i \mathcal{N}_2}{16\pi^2} \times \int_0^\infty d\ell_E^2 \frac{\ell_E^2}{(\ell_E^2 + \Delta)^3} \\ &= \frac{-i \mathcal{N}_2}{16\pi^2} \times \frac{1}{2\Delta} \\ &= \frac{-i}{32\pi^2} \times \frac{\mathcal{N}_2 = 4z(1-z)m^2 \quad \langle\langle \text{for } D = 4 \rangle\rangle}{\Delta = (1-z)^2 m^2 \quad \langle\langle \text{for } q^2 = 0 \rangle\rangle} \\ &= \frac{-i}{32\pi^2} \times \frac{4z}{1-z}. \end{aligned} \quad (54)$$

Substituting this formula into eq. (49), we have

$$F_2^{1\text{ loop}}(q^2 = 0) = \frac{e^2}{16\pi^2} \iiint_0^1 dx dy dz \delta(x + y + z - 1) \times \frac{4z}{1-z}. \quad (55)$$

The integrand here depends on z but not on the other two Feynman parameters, so we can immediately integrate over x and y and obtain

$$\iint_0^1 dx dy \delta(x + y + z - 1) = \int_0^{1-z} dx = 1 - z. \quad (56)$$

Consequently,

$$F_2^{1\text{loop}}(q^2 = 0) = \frac{e^2}{16\pi^2} \times \int_0^1 dz (1-z) \times \frac{4z}{1-z} = \frac{e^2}{16\pi^2} \times 2 = \frac{\alpha}{2\pi} \quad (57)$$

and the gyromagnetic moment is

$$g = 2 + \frac{\alpha}{\pi} + O(\alpha^2). \quad (58)$$

The higher-loop corrections to this gyromagnetic moment are harder to calculate because the number of diagrams grows very rapidly with the number of loops; at the 4-loop order there are thousands of diagrams, and one needs a computer just to count them! Also, at higher orders one has to include the effects of strong and weak interactions because the photons interact with hadrons and W^\pm particles, which in turn interact with other hadrons, Z^0 , Higgs, *etc.*, *etc.* Nevertheless, people have calculated the electron's and the muon's *anomalous magnetic moments*

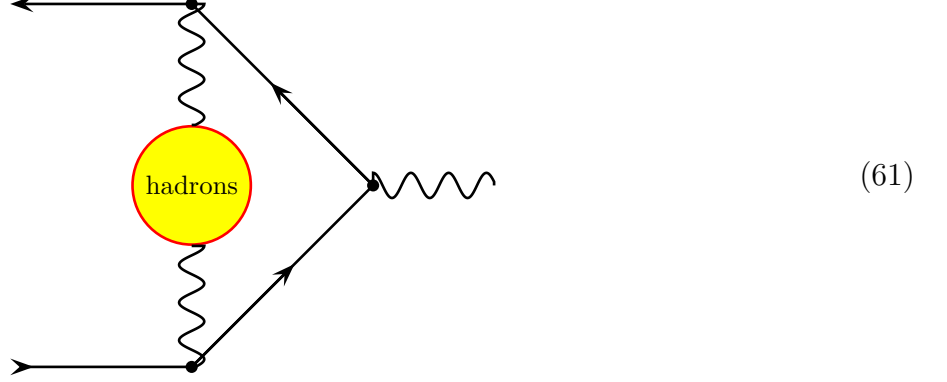
$$a_e = \frac{g_e - 2}{2} = F_2^{\text{electron}}(0) \quad \text{and} \quad a_\mu = \frac{g_\mu - 2}{2} = F_2^{\text{muon}}(0) \quad (59)$$

to the order α^4 back in the 1970s, and more recent calculations are good up to the order α^5 . Meanwhile, the experimentalists have measured a_e to a comparable accuracy of 12 significant digits and a_μ to 9 significant digits

$$a_e^{\text{exp}} = 0.001\,159\,652\,180\,62\,(12), \quad a_\mu^{\text{exp}} = 0.001\,165\,920\,59\,(22). \quad (60)$$

The theoretical value of the electron's anomalous magnetic moment is in good agreement with the experimental value, while for the muon there is a small discrepancy $a_\mu^{\text{exp}} - a_\mu^{\text{theory}} \approx (1.13 \pm 0.35) \cdot 10^{-9}$. This discrepancy might stem from some physics beyond the Standard Model, maybe supersymmetry, maybe something else. Note that the effect of heavy particles on the a_μ is proportional to $(m_\mu/M_{\text{heavy}})^2$, that's why the muon's anomalous magnetic moment is much more sensitive to the new physics than the electron's.

However, the discrepancy between the a_μ^{exp} and the a_μ^{theory} might also stem from a small theoretical error in modeling the photon-hadron interactions, which affects the a_μ^{theory} via 2+ loop diagrams like



For a recent review of the muon's high-precision anomalous magnetic moment — both the experiments and the theory — see [2024 review of the Particle Data Group, §10.4.2](#), and the references cited therein.

I would like to complete this section of the notes by calculating the $F_2^{1\text{loop}}(q^2)$ form factor for $q^2 \neq 0$. Proceeding as in eq. (54) but letting $\Delta = (1 - z)^2 m^2 - xyq^2$, we have

$$\int \frac{d^4 \ell}{(2\pi)^4} \frac{\mathcal{N}_2}{[\ell^2 - \Delta + i0]^3} = \frac{-i}{32\pi^2} \times \frac{4z(1-z)m^2}{(1-z)^2 m^2 - xyq^2} \quad (62)$$

and hence

$$F_2^{1\text{loop}}(q^2) = \frac{e^2}{16\pi^2} \iiint_0^1 dx dy dz \delta(x + y + z - 1) \times \frac{4z(1-z)m^2}{(1-z)^2 m^2 - xyq^2}. \quad (63)$$

To evaluate this integral over Feynman parameters, we change variables from x, y, z to $w = 1 - z$ and $\xi = x/(x + y)$,

$$\begin{aligned} x &= w\xi, & y &= w(1 - \xi), & z &= 1 - w, \\ dx dy dz \delta(x + y + z - 1) &= dx dy = w dw d\xi. \end{aligned} \quad (64)$$

Consequently,

$$\begin{aligned}
F_2^{1\text{loop}}(q^2) &= \frac{e^2}{16\pi^2} \int_0^1 d\xi \int_0^1 dw w \times \frac{4(1-w)w \times m^2}{w^2 \times m^2 - w^2\xi(1-\xi) \times q^2} \\
&= \frac{e^2}{16\pi^2} \int_0^1 d\xi \frac{m^2}{m^2 - \xi(1-\xi)q^2} \times \int_0^1 dw w \times \frac{4w(1-w)}{w^2} \\
&= \frac{e^2}{8\pi^2} \times \int_0^1 d\xi \frac{m^2}{m^2 - \xi(1-\xi)q^2} \\
&= \frac{\alpha}{2\pi} \times \frac{4m^2}{\sqrt{(-q^2) \times (4m^2 - q^2)}} \times \log \frac{\sqrt{4m^2 - q^2} + \sqrt{-q^2}}{2m} \\
&\quad \langle\langle \text{for } q^2 < 0 \rangle\rangle \\
&= \frac{\alpha}{2\pi} \times \frac{4m^2}{\sqrt{q^2 \times (4m^2 - q^2)}} \times \arctan \sqrt{\frac{q^2}{4m^2 - q^2}} \\
&\quad \langle\langle \text{for } 0 < q^2 < 4m^2 \rangle\rangle.
\end{aligned} \tag{65}$$

In particular, for large and spacelike q — *i.e.*, $q^2 < 0$ and $-q^2 \gg m^2$, —

$$F_2^{1\text{loop}}(q^2) \approx \frac{\alpha}{2\pi} \times \frac{2m^2}{-q^2} \times \log \frac{-q^2}{m^2}. \tag{66}$$

The Electric Form Factor

Now consider the electric form factor $F_1(q^2)$. In [the first section](#) we have obtained

$$F_1^{1\text{loop}}(q^2) = -2ie^2 \iiint_0^1 dx dy dz \delta(x+y+z-1) \int \frac{d^D \ell}{(2\pi)^D} \frac{\mathcal{N}_1}{[\ell^2 - \Delta + i0]^3}, \tag{48}$$

for

$$\mathcal{N}_1 \cong \frac{(D-2)^2}{D} \times \ell^2 - (D-2) \times \Delta + 2z \times (2m^2 - q^2) \tag{45}$$

and $\Delta = (1-z)^2 m^2 - xyq^2$.

Let's start by calculating the momentum integral in eq. (48). The numerator \mathcal{N}_1 depends on ℓ as $a\ell^2 + b$, so there is a logarithmic UV divergence for $\ell \rightarrow \infty$; to regularize this divergence, we work in $D = 4 - 2\epsilon$ dimensions. Thus,

$$\begin{aligned}
-i \int_{\text{reg}} \frac{d^4 \ell}{(2\pi)^4} \frac{a\ell^2 + b}{[\ell^2 - \Delta + i0]^3} &\equiv -i\mu^{4-D} \int \frac{d^D \ell}{(2\pi)^D} \frac{a\ell^2 + b}{[\ell^2 - \Delta + i0]^3} = \\
&= -i\mu^{4-D} \int \frac{id^D \ell_E}{(2\pi)^D} \frac{-a\ell_E^2 + b}{-[\ell_E^2 + \Delta]^3} \\
&= \mu^{4-D} \int \frac{d^D \ell_E}{(2\pi)^D} \times \left[\frac{a\ell_E^2 - b}{(\ell_E^2 + \Delta)^3} = \frac{a}{(\ell_E + \Delta)^2} - \frac{a\Delta + b}{(\ell_E^2 + \Delta)^3} \right] \\
&= \mu^{4-D} \int \frac{d^D \ell_E}{(2\pi)^D} \int_0^\infty dt \left(a \times t - (a\Delta + b) \times \frac{1}{2}t^2 \right) \times e^{-t(\Delta + \ell_E^2)} \\
&= \int_0^\infty dt \left(a \times t - (a\Delta + b) \times \frac{1}{2}t^2 \right) e^{-t\Delta} \times \mu^{4-D} \int \frac{d^D \ell_E}{(2\pi)^D} e^{-t\ell_E^2} \\
&= \int_0^\infty dt \left(a \times t - (a\Delta + b) \times \frac{1}{2}t^2 \right) e^{-t\Delta} \times \frac{\mu^{4-D}}{(4\pi t)^{D/2}} \\
&= \frac{\mu^{4-D}}{(4\pi)^{D/2}} \int_0^\infty dt e^{-t\Delta} \times \left(a \times t^{1-(D/2)} - \frac{1}{2}(a\Delta + b) \times t^{2-(D/2)} \right) \\
&= \frac{\mu^{4-D}}{(4\pi)^{D/2}} \left\{ a \times \Gamma\left(2 - \frac{D}{2}\right) \times \Delta^{\frac{D}{2}-2} - \frac{1}{2}(a\Delta + b) \times \Gamma\left(3 - \frac{D}{2}\right) \times \Delta^{\frac{D}{2}-3} \right\} \\
&= \frac{\mu^{4-D}}{(4\pi)^{D/2}} \Gamma\left(2 - \frac{D}{2}\right) \Delta^{\frac{D}{2}-2} \times \left\{ a - \frac{1}{2}(a\Delta + b) \times \frac{2 - \frac{D}{2}}{\Delta} \right\} \\
&\rightarrow \frac{(4\pi\mu^2)^\epsilon}{16\pi^2} \times \frac{\Gamma(\epsilon)}{\Delta^\epsilon} \times \left\{ a - \frac{\epsilon}{2} \times \frac{a\Delta + b}{\Delta} \right\}.
\end{aligned} \tag{67}$$

Going back to eq. (45), we identify a and b in eq. (66) as

$$\begin{aligned}
a &= \frac{(D-2)^2}{D} = \frac{2(1-\epsilon)^2}{2-\epsilon}, \\
b &= 2z \times (2m^2 - q^2) - (D-2) \times \Delta = \hat{b} - 2(1-\epsilon)\Delta,
\end{aligned} \tag{68}$$

where $\hat{b} = 2z(2m^2 - q^2)$.

Consequently, on the last line of eq. (67) we have

$$\begin{aligned}
\{\dots\} &= a - \frac{\epsilon}{2} \times \frac{a\Delta + b = a\Delta + \hat{b} - 2(1-\epsilon)\Delta}{\Delta} \\
&= \left(a = \frac{2(1-\epsilon)^2}{2-\epsilon} \right) \times \left(1 - \frac{\epsilon}{2} \right) - \frac{\epsilon}{2} \times \frac{\hat{b}}{\Delta} + \frac{\epsilon}{2} \times 2(1-\epsilon) \\
&= (1-\epsilon)^2 + \epsilon(1-\epsilon) - \frac{\epsilon}{2} \times \frac{\hat{b}}{\Delta} \\
&= (1-\epsilon) - \frac{\epsilon}{2} \times \frac{\hat{b}}{\Delta},
\end{aligned} \tag{69}$$

so the momentum integral for the electric form factors evaluates to

$$-2ie^2 \mu^{4-D} \int \frac{d^D \ell}{(2\pi)^D} \frac{\mathcal{N}_1}{[\ell^2 - \Delta + i0]^3} = \frac{2e^2}{16\pi^2} \Gamma(\epsilon) \left(\frac{4\pi\mu^2}{\Delta} \right)^\epsilon \left\{ (1-\epsilon) - \frac{\epsilon}{2} \times \frac{\hat{b}}{\Delta} \right\}. \tag{70}$$

The next step in our calculation is to integrate the result in eq. (70) over the Feynman parameters. Changing the integration variables from x, y, z to w and ξ according to

$$\begin{aligned}
x &= w\xi, \quad y = w(1-\xi), \quad z = 1-w, \\
dx dy dz \delta(x+y+z-1) &= dx dy = w dw d\xi,
\end{aligned} \tag{64}$$

we get

$$\hat{b} = (1-w) \times (2m^2 - q^2) \tag{71}$$

and

$$\Delta(w, \xi) = (1-z)^2 m^2 - xyq^2 = w^2 \times (m^2 - \xi(1-\xi)q^2), \tag{72}$$

or equivalently

$$\Delta(w, \xi) = w^2 \times H(\xi) \quad \text{where} \quad H(\xi) \stackrel{\text{def}}{=} m^2 - \xi(1-\xi)q^2. \tag{73}$$

Altogether, this gives us

$$\begin{aligned}
F_1^{\text{1loop}}(q^2) &= \frac{\alpha}{2\pi} \Gamma(\epsilon) (4\pi\mu^2)^\epsilon \times \int_0^1 d\xi \int_0^1 dw w \times \left\{ \frac{(1-\epsilon)}{[\Delta(w, \xi)]^\epsilon} - \epsilon \frac{(1-w)(2m^2 - q^2)}{[\Delta(w, \xi)]^{1+\epsilon}} \right\} \\
&= \frac{\alpha}{2\pi} \Gamma(\epsilon) (4\pi\mu^2)^\epsilon \times \int_0^1 d\xi \int_0^1 dw \left\{ \frac{(1-\epsilon)}{H^\epsilon} \times \frac{w}{w^{2\epsilon}} - \epsilon \times \frac{2m^2 - q^2}{H^{1+\epsilon}} \times \frac{w(1-w)}{w^{2+2\epsilon}} \right\}.
\end{aligned} \tag{74}$$

In this double integral we $\int dw$ before we $\int d\xi$, and the first stage has an immediate trouble in the $w \rightarrow 0$ limit of the w integral: the second term inside $\{\dots\}$ behaves as $w^{-1-2\epsilon}$, which is not integrable for $\epsilon \geq 0$ — *i.e.*, for any dimension $D \leq 4$. Indeed,

$$\int_0^{\text{positive}} \frac{dw}{w^{1+2\epsilon}} = \infty \quad \text{for } \epsilon \geq 0 : \quad (75)$$

For $\epsilon = 0$ ($D = 4$) the divergence here is logarithmic while for $\epsilon > 0$ ($D < 4$) it becomes power-like.

The Infrared Divergence

Physically, the divergence (75) is *infrared* rather than ultraviolet, that's why it gets worse as we lower the dimension D . Indeed, let's go back to the diagram (9) and look at the denominator \mathcal{D} in eqs. (10) and (12). Taking the electron's momenta p and p' on-shell before introducing the Feynman parameters, we have

$$(p+k)^2 - m^2 = k^2 + 2kp + \cancel{p^2} - \cancel{m^2} = k^2 + 2kp = O(|k|) \quad \text{when } k \rightarrow 0, \quad (76)$$

and likewise

$$(p'+k)^2 - m^2 = k^2 + 2kp' = O(|k|) \quad \text{when } k \rightarrow 0. \quad (77)$$

Combining these two electron propagators with the $O(1/k^2)$ photon propagators, we see that the net denominator behaves as $\mathcal{D} \propto |k|^4$ for $k \rightarrow 0$. At the same time, the numerator \mathcal{N}^μ remains finite, which makes the integral

$$\int d^D k \frac{\mathcal{N}^\mu}{\mathcal{D}} \propto \int \frac{d^D k}{|k|^4} \quad (78)$$

diverge for $k \rightarrow 0$. In $D = 4$ dimensions, the infrared divergence here is logarithmic, while in lower dimensions $D < 4$ it becomes power-like, *i.e.* $O((1/k_{\min})^{4-D})$ — precisely as in eqs. (75) and (74).

We can regularize the infrared divergence (78) — and also (75) — by analytically continuing the spacetime dimension to $D > 4$. Such dimensional regularization of the IR divergences is used in many situations in both QFT and condensed matter. However, taking $D > 4$ makes the ultraviolet

divergences worse, so if some amplitude has both UV and IR divergences, we cannot cure both of them at the same time by analytically continuing to $D \neq 4$. In particular, when calculating the electric form factor $F_1(q^2)$ of the electron, we need $D < 4$ to regulate the momentum integral $\int d^D \ell$, but then we need $D > 4$ to regulate the integral over the Feynman parameters.

A common *dirty trick* is to first continue to $D < 4$, shift the loop momentum from k^μ to $\ell^\mu = k^\mu + \text{shift}$, evaluate the $\int d^D \ell$ momentum integral in $D < 4$ dimension, then analytically continue the result to $D > 4$ and integrate over the Feynman parameters, and ultimately continue the final result to $D = 4$. However, in this kind of dimensional regularization it is hard to disentangle the $1/\epsilon$ poles coming from the UV divergence $\log(\Lambda^2/\mu^2)$ from the $1/\epsilon$ poles coming from the IR divergence $\log(\mu^2/k_{\min}^2)$, so we are not going to use it here.

Instead, we are going to use DR for the UV divergence only, while the IR divergence is regulated by a tiny but not-quite-zero photon mass $m_\gamma^2 \ll m_e^2$. Strictly speaking, a massive vector particle has three polarization states and its propagator is

$$\text{wavy line} = \frac{-i}{k^2 - m_\gamma^2 + i0} \times \left(g^{\mu\nu} - \frac{k^\mu k^\nu}{m_\gamma^2} \right). \quad (79)$$

However, if we assume that the mass is due to Higgs mechanism and use the R_ξ gauge for $\xi = 1$ — see [my notes on the Higgs mechanism, pages 6–8](#) for details, — then the propagator becomes simply Feynman-like

$$\text{wavy line} = \frac{-ig^{\mu\nu}}{k^2 - m_\gamma^2 + i0}. \quad (80)$$

The price of an R_ξ gauge is the persistence of the Goldstone scalar — which should be eaten by the Higgs mechanism — as a scalar field with $\text{mass}_g^2 = \xi m_\gamma^2$. However, this unphysical scalar field would not affect our calculations at this level of analysis.

Using this infrared regulator for the internal photon line in the one-loop diagram (9), we get the vertex amplitude that looks exactly like eq. (10) except for one factor in the denominator,

$$\frac{1}{k^2 + i0} \text{ becomes } \frac{1}{k^2 - m_\gamma^2 + i0}. \quad (81)$$

In terms of the integral (22), this change has no effect on the numerator \mathcal{N}^μ , but the denominator

becomes $[\ell^2 - \tilde{\Delta} + i0]^3$ where

$$\ell^2 - \tilde{\Delta}(x, y, z) = x \times ((p+k)^2 - m^2) + y \times ((p'+k)^2 - m^2) + z \times (k^2 - m_\gamma^2) \quad (82)$$

hence

$$\ell = k + xp + yp' \quad (83)$$

exactly as in eq. (20) but

$$\tilde{\Delta}(x, y, z) = (1-z)^2 m^2 - xyq^2 + z \times m_\gamma^2 = \Delta(x, y, z) + z \times m_\gamma^2. \quad (84)$$

Consequently, the electric form factor is

$$F_1^{\text{1loop}}(q^2) = \int d(FP) \int \mu^{4-D} \frac{d^D \ell}{(2\pi)^D} \frac{-2ie^2 \times \mathcal{N}_1}{[\ell^2 - \tilde{\Delta} + i0]^3}, \quad (85)$$

exactly as in eq. (48), except for the $\tilde{\Delta}$ instead of the original Δ in the denominator. The momentum integral here converges for any $D < 4$ and it evaluates exactly as in eq. (67). The only subtlety here is that in the numerator, the ℓ -independent term b involves the un-modified Δ instead of $\tilde{\Delta}$ (*cf.* eq. (68)), but we can fix that by redefining

$$\hat{b} \rightarrow \tilde{b} = 2z(2m_e^2 - q^2 + (1-\epsilon)m_\gamma^2) \implies \text{same old } b = \tilde{b} - 2(1-\epsilon)\tilde{\Delta} \quad (86)$$

Consequently, the momentum integral in $D < 4$ dimensions evaluates to

$$\frac{\alpha}{2\pi} \Gamma(\epsilon) \left(\frac{4\pi\mu^2}{\tilde{\Delta}} \right)^\epsilon \left\{ (1-\epsilon) - \frac{\epsilon}{2} \times \frac{\tilde{b}}{\tilde{\Delta}} \right\}, \quad (87)$$

and hence the IR-regulated F_1 form-factor becomes

$$F_1^{\text{1loop}}(q^2) = \frac{\alpha}{2\pi} \Gamma(\epsilon) (4\pi\mu^2)^\epsilon \int_0^1 d\xi \int_0^1 dw w \times \left\{ \frac{1-\epsilon}{[\tilde{\Delta}(w, \xi)]^\epsilon} - \frac{\epsilon(1-w)(2m_e^2 - q^2 + (1-\epsilon)m_\gamma^2)}{[\tilde{\Delta}(w, \xi)]^{1+\epsilon}} \right\} \quad (88)$$

where

$$\tilde{\Delta}(w, \xi) = (1-z)^2 m_e^2 - xyq^2 + zm_\gamma^2 = w^2 \times H(\xi) + (1-w) \times m_\gamma^2. \quad (89)$$

Note that the photon's mass is tiny, $m_\gamma^2 \ll m_e^2, q^2$; were it not for the IR divergences, we would have used $m_\gamma^2 = 0$. This allows us to neglect various $O(m_\gamma^2)$ terms in eq. (88) except when

it would cause a divergence for $w \rightarrow 0$; in particular, we may neglect the $(1 - \epsilon)m_\gamma^2$ term in the numerator of the second term in the integrand. As to the denominators, in eq. (89) the second term containing the photon's mass becomes important only in the $w \rightarrow 0$ limit, and in that limit $(1 - w)m_\gamma^2 \rightarrow m_\gamma^2$. Thus, we approximate

$$\tilde{\Delta}(w, \xi) \approx w^2 \times H(\xi) + m_\gamma^2 \quad (90)$$

and the $\int dw$ integral in eq. (88) becomes

$$\begin{aligned} \int_0^1 dw w \times \left\{ \frac{1 - \epsilon}{[w^2 H(\xi) + m_\gamma^2]^\epsilon} - \frac{\epsilon(1 - w)(2m_e^2 - q^2)}{[w^2 H(\xi) + m_\gamma^2]^{1+\epsilon}} \right\} \\ = \frac{(1 - \epsilon)}{H^\epsilon} \times \int_0^1 \frac{dw w}{[w^2 + (m_\gamma^2/H)]^\epsilon} \\ + \frac{\epsilon(2m_e^2 - q^2)}{H^{1+\epsilon}} \times \int_0^1 \frac{dw w^2}{[w^2 + (m_\gamma^2/H)]^{1+\epsilon}} \\ - \frac{\epsilon(2m_e^2 - q^2)}{H^{1+\epsilon}} \times \int_0^1 \frac{dw w}{[w^2 + (m_\gamma^2/H)]^{1+\epsilon}}. \end{aligned} \quad (91)$$

For $0 < \epsilon < \frac{1}{2}$ — *i.e.*, for $3 < D < 4$ — the integrals on the second and third lines here converge even for $m_\gamma^2 = 0$,

$$\begin{aligned} \int_0^1 \frac{dw w}{[w^2]^\epsilon} &= \frac{1}{2 - 2\epsilon} \quad \text{for } \epsilon < 1, \\ \int_0^1 \frac{dw w^2}{[w^2]^{1+\epsilon}} &= \frac{1}{1 - 2\epsilon} \quad \text{for } \epsilon < \frac{1}{2}, \end{aligned} \quad (92)$$

so we may just as well evaluate them without the photon's mass. Only on the last line of eq. (91) we do need $m_\gamma^2 \neq 0$ to make the integral converge for some $D \leq 4$:

$$\int_0^1 \frac{dw w}{[w^2 + (m_\gamma^2/H)]^{1+\epsilon}} = \frac{-1}{2\epsilon} \frac{1}{[w^2 + (m_\gamma^2/H)]^\epsilon} \Big|_0^1 = \frac{1}{2\epsilon} \left[\left(\frac{H}{m_\gamma^2} \right)^\epsilon - 1 \right]. \quad (93)$$

Combining all these $\int dw$ integrals together, we get

$$\begin{aligned}
\int_0^1 dw \{ \dots \} &= \frac{(1-\epsilon)}{H^\epsilon} \times \frac{1}{2(1-\epsilon)} + \frac{\epsilon(2m_e^2 - q^2)}{H^{1+\epsilon}} \times \frac{1}{1-2\epsilon} \\
&\quad - \frac{\epsilon(2m_e^2 - q^2)}{H^{1+\epsilon}} \times \frac{1}{2\epsilon} \left[\left(\frac{H}{m_\gamma^2} \right)^\epsilon - 1 \right] \\
&= \frac{1}{2H^\epsilon} + \frac{2m^2 - q^2}{2H^{1+\epsilon}} \times \left[\frac{2\epsilon}{1-2\epsilon} - \left(\frac{H}{m_\gamma^2} \right)^\epsilon + 1 \right] \\
&= \frac{1}{2H^\epsilon} + \frac{2m^2 - q^2}{2H^{1+\epsilon}} \times \left[\frac{1}{1-2\epsilon} - \left(\frac{H}{m_\gamma^2} \right)^\epsilon \right]
\end{aligned} \tag{94}$$

and hence

$$F_1^{1\text{loop}}(q^2) = \frac{\alpha}{4\pi} \int_0^1 d\xi \Gamma(\epsilon) \left(\frac{4\pi\mu^2}{H(\xi)} \right)^\epsilon \times \left\{ 1 + \frac{2m_e^2 - q^2}{H(\xi)} \times \left[\frac{1}{1-2\epsilon} - \left(\frac{H(\xi)}{m_\gamma^2} \right)^\epsilon \right] \right\} \tag{95}$$

where

$$H(\xi) = m_e^2 - \xi(1-\xi)q^2. \tag{73}$$

Before we even try to perform this last integral, let's remember that

$$\Gamma_{\text{net}}^\mu = \gamma_{\text{tree}}^\mu + \Gamma_{\text{loops}}^\mu + \delta_1 \times \gamma^\mu \tag{5}$$

and hence

$$F_1^{\text{net}}(q^2) = 1^{\text{tree}} + F_1^{\text{loops}}(q^2) + \delta_1. \tag{51}$$

Also, the net electric charge does not renormalize, so we must have

$$F_1^{\text{net}}(q^2) \rightarrow 1 \quad \text{for } q^2 \rightarrow 0 \quad \implies \quad \delta_1 = -F_1^{\text{loops}}(q^2 = 0). \tag{8}$$

In particular, the δ_1 counterterm to the order α follows from eq. (95) for $q^2 = 0$, in which case $H(\xi) \equiv m_e^2$ and the $\int d\xi$ integral becomes trivial (the integrand does not depend on ξ at all).

Thus,

$$\delta_1 = -\frac{\alpha}{4\pi} \Gamma(\epsilon) \left(\frac{4\pi\mu^2}{m_e^2} \right)^\epsilon \times \left\{ 1 + \frac{2}{1-2\epsilon} - 2 \left(\frac{m_e^2}{m_\gamma^2} \right)^\epsilon \right\} + O(\alpha^2). \quad (96)$$

This formula holds for any dimension D between 3 and 4 (*i.e.*, $0 < \epsilon < \frac{1}{2}$). In the $D \rightarrow 4$ limit, it becomes

$$\delta_1 = -\frac{\alpha}{4\pi} \times \left\{ \frac{1}{\epsilon} - \gamma_E + \log \frac{4\pi\mu^2}{m_e^2} + 4 - 2 \log \frac{m_e^2}{m_\gamma^2} \right\} + O(\alpha^2). \quad (97)$$

Now let's go back to the electric form factor $F_1^{\text{net}}(q^2)$ for $q^2 \neq 0$. According to eqs. (51) and (8), at the one-loop level

$$F_1^{\text{net}}(q^2) - 1 = F_1^{1\text{loop}}(q^2) - F_1^{1\text{loop}}(0) + O(\alpha^2) \quad (98)$$

where $F_1^{1\text{loop}}(q^2)$ is given by eq. (95). Taking the $\epsilon \rightarrow 0$ limit of that formula, we arrive at

$$F_1^{1\text{loop}}(q^2) = \frac{\alpha}{4\pi} \int_0^1 d\xi \left\{ \frac{1}{\epsilon} - \gamma_E + \log \frac{4\pi\mu^2}{H(\xi)} + \frac{2m_e^2 - q^2}{H(\xi)} \times \left[2 - \log \frac{H(\xi)}{m_\gamma^2} \right] \right\}, \quad (99)$$

and now we should subtract a similar expression for $q^2 = 0$. This subtraction cancels the UV divergence and the associated $1/\epsilon$ pole but not the IR divergence. Moreover, not only the subtracted one-loop amplitude depends on the IR regulators, but the coefficient of the $\log m_\gamma^2$ has a non-trivial momentum dependence. Indeed,

$$\begin{aligned} F_1^{1\text{loop}}(q^2) - F_1^{1\text{loop}}(0) &= \\ &= \frac{\alpha}{4\pi} \int_0^1 d\xi \left\{ \begin{array}{l} \frac{1}{\epsilon} - \gamma_E + \log \frac{4\pi\mu^2}{H(\xi)} + \frac{2m_e^2 - q^2}{H(\xi)} \times \left[2 - \log \frac{H(\xi)}{m_\gamma^2} \right] \\ -\frac{1}{\epsilon} + \gamma_E - \log \frac{4\pi\mu^2}{m_e^2} - 2 \times \left[2 - \log \frac{m_e^2}{m_\gamma^2} \right] \end{array} \right\} \\ &= \frac{\alpha}{4\pi} \int_0^1 d\xi \left\{ \log \frac{m_e^2}{H(\xi)} + \frac{2m_e^2 - q^2}{H(\xi)} \times \left[2 - \log \frac{m_e^2}{m_\gamma^2} - \log \frac{H(\xi)}{m_e^2} \right] - 2 \left[2 - \log \frac{m_e^2}{m_\gamma^2} \right] \right\} \end{aligned}$$

$$\begin{aligned}
&= \frac{\alpha}{4\pi} \left[2 - \log \frac{m_e^2}{m_\gamma^2} \right] \times \int_0^1 d\xi \left(\frac{2m_e^2 - q^2}{H(\xi)} - 2 \right) \\
&\quad + \frac{\alpha}{4\pi} \int_0^1 d\xi \log \frac{m_e^2}{H(\xi)} \times \left(1 + \frac{2m_e^2 - q^2}{H(\xi)} \right). \tag{100}
\end{aligned}$$

Following the textbook notations, let's explicitly separate the IR divergent part here from the finite part, thus

$$\begin{aligned}
F_1^{\text{order } \alpha}(q^2) &= 1 - \frac{\alpha}{2\pi} \times \left\{ f_{\text{IR}}(q^2) \times \log \frac{m_e^2}{m_\gamma^2} + f_{\text{fin}}(q^2) \right\} + O(\alpha^2) \\
&= 1 - \frac{\alpha}{2\pi} \times \left\{ f_{\text{IR}}(q^2) \times \log \frac{-q^2}{m_\gamma^2} + \tilde{f}_{\text{fin}}(q^2) \right\} + O(\alpha^2) \tag{101}
\end{aligned}$$

where

$$f_{\text{IR}}(q^2) = \int_0^1 d\xi \left(\frac{2m_e^2 - q^2}{2H(\xi)} - 1 \right), \tag{102}$$

$$f_{\text{fin}}(q^2) = \frac{1}{2} \int_0^1 d\xi \left(\frac{2m_e^2 - q^2}{H(\xi)} + 1 \right) \times \log \frac{H(\xi)}{m_e^2} - 2f_{\text{IR}}(q^2), \tag{103}$$

$$\tilde{f}_{\text{fin}}(q^2) = f_{\text{fin}}(q^2) - f_{\text{IR}}(q^2) \times \log \frac{-q^2}{m_e^2}. \tag{104}$$

All three integrals here vanish for $q^2 = 0$ — which upholds the $F_1(0) = 1$ requirement — but have finite non-zero values for all other q^2 . In particular, $f_{\text{IR}}(q^2) \neq 0$, which means that the $F_1(q^2)$ form factor (101) suffers from IR divergence at all $q^2 \neq 0$.

Also, for very large q^2 , the $F_1^{\text{order } \alpha} - 1$ grows like a quadratic polynomial in $\log(-q^2/m^2)$. Indeed, in this limit

$$f_{\text{IR}}(q^2) \approx 2 \log \frac{-q^2}{m_e^2} - 1, \quad f_{\text{fin}}(q^2) \approx \frac{1}{2} \log^2 \frac{-q^2}{m_e^2} - \frac{7}{2} \log \frac{-q^2}{m_e^2} + 1 - \frac{\pi^2}{6}, \tag{105}$$

hence

$$F_1(q^2) \approx 1 - \frac{\alpha}{2\pi} \left\{ \frac{1}{2} \log^2 \frac{-q^2}{m_e^2} + \left(2 \log \frac{m_e^2}{m_\gamma^2} - \frac{7}{2} \right) \log \frac{-q^2}{m_e^2} - 2 \log \frac{m_e^2}{m_\gamma^2} + \text{const} \right\} + O(\alpha^2). \tag{106}$$

The leading \log^2 term here is known as the *Sudakov's double logarithm*; it plays important role

in estimating radiative corrections (*i.e.*, loop corrections) to various QED processes involving relativistic electrons. Unfortunately, I do not have class time to discuss it in any detail, but if you are interested, you may look it up in [this Lipatov’s review](#). Briefly, for any process involving high-energy electron scattering, electron-positron collisions, or leptonic decays of heavy particles, the relative magnitude of QED correction is

$$\text{correction} \sim \frac{\alpha}{\pi} \times \log^2 \frac{E_e}{m_e}, \quad (107)$$

which becomes surprisingly large for $E_e \gtrsim 100$ GeV. And if we ever accelerate electrons to multi-TeV energies, or make particles with multi-TeV masses and observe their decays into electrons (plus some other particles), we would not be able to calculate those processes in perturbation theory because L -loop QED corrections due to Sudakov’s double logarithms become

$$O\left(\left(\frac{\alpha}{\pi} \log^2 \frac{-q^2}{m_e^2}\right)^L\right) = O(1). \quad (108)$$

Meanwhile, let’s turn our attention to a more urgent question: How does the infrared divergence of the F_1 form factor affect the physical cross-sections of various QED processes? In the [next section](#) I shall explain this issue in some detail, but let me summarize the bottom line here:

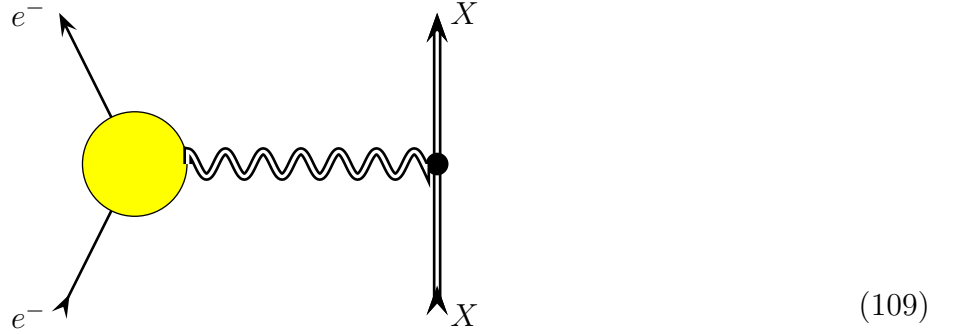
- ★ The S-matrix elements in QED between specific n -particle states are infrared-divergent, so the *exclusive* scattering cross-sections such as $\sigma(e^- + X \rightarrow e^- + X + \text{nothing else})$ are all infrared-divergent.
- ★ However, the more inclusive cross-sections such as $\sigma(e^- + X \rightarrow e^- + X + \text{optional soft photons})$ are perfectly finite: all the IR divergences cancel out. By *optional* soft photons I mean photons which would not be detected in a real-life particle experiment because their energies are too low, so we do not know if they are present or absent in the final state.

Finite Cross-sections for IR–Divergent Amplitudes

In the [previous section](#) we saw that the one-loop correction to the electron’s electric form factor $F_1(q^2)$ suffers from an infrared divergence. Many other QED amplitudes — essentially, all the amplitudes involving on-shell electrons or positrons — also suffer from similar IR divergences at the one-loop and higher-loop levels. At the same time, the processes involving emission of

soft (*i.e.*, low-energy) photons have IR-divergent cross-sections already at the tree level; this is explained in detail in §6.1 of the Peskin&Schroeder textbook. But somehow, the two kinds of infrared divergences cancel out from the inclusive cross-sections in which we allow for optional extra soft photons in the final state.

For example, consider elastic Coulomb scattering of an electron off a very heavy point-like particle X :



For $q^2 \ll M_X^2$, we may approximate the X particle as a static (non-recoiling) source of electric field, and in its rest frame (the lab frame), the scattering amplitude evaluates to

$$\frac{\mathcal{M}(eX \rightarrow eX)}{2M_X} = \frac{4\pi\alpha_{\text{eff}}(q^2)}{q^2} \times \bar{u}(p')\Gamma^0(p', p)u(p). \quad (110)$$

At the tree level $\Gamma^\mu(p', p) \equiv \gamma^\mu$, while at the one-loop level we get non-trivial form-factors. Focusing at their infrared divergences, we have

$$F_2^{\text{net}}(q^2) = \frac{\alpha}{2\pi} \times \text{finite} + O(\alpha^2) \quad (111)$$

while

$$F_1^{\text{net}}(q^2) = 1 - \frac{\alpha}{2\pi} \left(f_{\text{IR}}(q^2) \times \log \frac{m_e^2}{m_\gamma^2} + \text{finite} \right) + O(\alpha^2) \quad (112)$$

and therefore

$$\Gamma_{\text{order } \alpha}^\mu(p', p) = \left(1 - \frac{\alpha}{2\pi} \times f_{\text{IR}}(q^2) \times \log \frac{m_e^2}{m_\gamma^2} \right) \times \Gamma_{\text{tree}}^\mu + \frac{\alpha}{2\pi} \times \text{finite}. \quad (113)$$

Plugging this one-loop vertex into the Coulomb scattering amplitude (110), we obtain

$$\mathcal{M}^{\text{tree}+1\text{ loop}}(eX \rightarrow eX) = \mathcal{M}^{\text{tree}}(eX \rightarrow eX) \times \left(1 - \frac{\alpha}{2\pi} \times f_{\text{IR}}(q^2) \times \log \frac{m_e^2}{m_\gamma^2} + \frac{\alpha}{2\pi} \times \text{finite} \right) \quad (114)$$

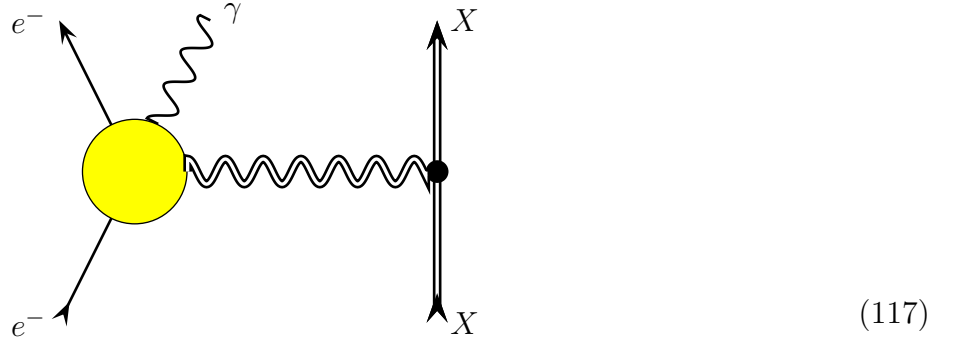
and hence the partial cross-section

$$\frac{d\sigma^{\text{tree}+1\text{ loop}}(eX \rightarrow eX)}{d\Omega} = \frac{d\sigma^{\text{tree}}(eX \rightarrow eX)}{d\Omega} \times \left(1 - \frac{\alpha}{\pi} \times f_{\text{IR}}(q^2) \times \log \frac{m_e^2}{m_\gamma^2} + \frac{\alpha}{\pi} \times \text{finite} + O(\alpha^2)\right). \quad (115)$$

Note the IR divergence of the one-loop term in this cross-section. For future reference, I would like to rephrase it in terms of $\log(E_e/m_\gamma)$ rather than $\log(m_e/m_\gamma)$, thus

$$\frac{d\sigma^{\text{tree}+1\text{ loop}}(eX \rightarrow eX)}{d\Omega} = \frac{d\sigma^{\text{tree}}(eX \rightarrow eX)}{d\Omega} \times \left(1 - \frac{\alpha}{\pi} \times f_{\text{IR}}(q^2) \times 2 \log \frac{E_e}{m_\gamma} + \frac{\alpha}{\pi} \times \text{finite} + O(\alpha^2)\right). \quad (116)$$

However, while the *exclusive* cross-section $\sigma(eX \rightarrow eX + \text{nothing else})$ suffer from the infrared divergence, the *inclusive* cross-sections $\sigma(eX \rightarrow eX + \text{optional photons})$ are perfectly finite. Let's see how this works at the order $O(\alpha) \times \sigma_{\text{tree}}(eX \rightarrow eX)$. In general, each additional incoming or outgoing photon costs a factor e in the amplitude and hence a factor α in the cross-section. In particular, the inelastic scattering in which a photon is emitted, $eX \rightarrow eX\gamma$,



has tree-level cross-section

$$\sigma_{\text{tree}}(eX \rightarrow eX\gamma) \sim \alpha \times \sigma_{\text{tree}}(eX \rightarrow eX) \sim \Delta_{1\text{ loop}}\sigma(eX \rightarrow eX), \quad (118)$$

so we may hope that the IR divergence of this tree-level cross-section would cancel the IR divergence of the one-loop exclusive cross-section (116).

To see how this cancellation works, let's focus on the *soft photon* regime of the inelastic scattering (117) in which the photon has much less energy than the incoming or the outgoing electron, $\omega_\gamma \ll E_e, E'_e$. In the [next section of these notes](#), we shall see that in this regime the partial cross-section integrated over the photon's directions but not its frequency or the electron's direction is

$$\frac{d\sigma^{\text{tree}}(eX \rightarrow eX\gamma)}{d\Omega_e d\omega_\gamma} = \frac{d\sigma^{\text{tree}}(eX \rightarrow eX)}{d\Omega_e} \times \frac{\alpha}{\pi} \times \frac{2f_{\text{IR}}(q^2)}{\omega_\gamma} \quad (119)$$

where the coefficient $f_{\text{IR}}(q^2)$ is precisely as in eq. (102) for the IR divergence of the F_1 form factor. Further integrating this partial cross-section over the photon frequencies ω_γ , we immediately run into the infrared divergence:

$$\int d\omega_\gamma \frac{d\sigma}{d\omega_\gamma} \propto \int_0^{\omega_{\text{max}} \sim E_e} \frac{d\omega_\gamma}{\omega_\gamma} = \log \frac{\omega_{\text{max}} \sim E_e}{\omega_{\text{min}} = 0} = \infty. \quad (120)$$

To regulate this divergence, we need to impose a minimal energy requirement on the emitted photon, and the simplest way to do this is to assume a tiny but non-zero photon mass m_γ . Consequently

$$\int_{\text{reg}} \frac{d\omega_\gamma}{\omega_\gamma} = \log \frac{\omega_{\text{max}} \sim E_e}{\omega_{\text{min}} \sim m_\gamma} + \text{finite} = \log \frac{E_e}{m_\gamma} + \text{finite}, \quad (121)$$

and therefore

$$\frac{d\sigma^{\text{tree}}(eX \rightarrow eX\gamma)}{d\Omega} = \frac{d\sigma^{\text{tree}}(eX \rightarrow eX)}{d\Omega} \times \frac{\alpha}{\pi} \left(2f_{\text{IR}} \times \log \frac{E_e}{m_\gamma} + \text{finite} \right). \quad (122)$$

Note that the IR divergence of this tree-level cross-section is precisely the same as of the one-loop cross-section (116), except for opposite signs,

$$\begin{aligned} \frac{d\sigma^{\text{tree}+1\text{ loop}}(Xe \rightarrow Xe)}{d\Omega} &= \frac{d\sigma^{\text{tree}}(eX \rightarrow eX)}{d\Omega} \times \\ &\times \left(1 - \frac{\alpha}{\pi} \times 2f_{\text{IR}}(q^2) \log \frac{E_e}{m_\gamma} + \frac{\alpha}{\pi} \times \text{finite} + O(\alpha^2) \right), \end{aligned} \quad (116)$$

$$\begin{aligned} \frac{d\sigma^{\text{tree}}(Xe \rightarrow Xe\gamma)}{d\Omega} &= \frac{d\sigma^{\text{tree}}(eX \rightarrow eX)}{d\Omega} \times \\ &\times \left(0 + \frac{\alpha}{\pi} \times 2f_{\text{IR}}(q^2) \log \frac{E_e}{m_\gamma} + \frac{\alpha}{\pi} \times \text{finite} \right), \end{aligned} \quad (122)$$

and therefore **the combined cross-section has no infrared divergence:**

$$\frac{d\sigma^{\text{tree}+1\text{ loop}}(Xe \rightarrow Xe) + d\sigma^{\text{tree}}(Xe \rightarrow Xe\gamma)}{d\Omega_e} = \frac{d\sigma^{\text{tree}}(Xe \rightarrow Xe)}{d\Omega_e} \times \left(1 + \frac{\alpha}{\pi} \times \text{finite} + O(\alpha^2)\right). \quad (123)$$

But what do we do with the IR divergences of the partial cross-sections (116) and (122)? While it is OK to UV-regulate or IR-regulate the intermediate stages of a calculation, the final result for a measurable quantity like a partial cross-section must be finite and it cannot depend on an IR regulator like m_γ . Nevertheless, eqs. (122) and (116) seem to contradict this rule, so what gives?

To resolve this paradox, consider a real-life scattering experiment. No photon detector can detect a photon with an arbitrarily low energy ω_γ , there is always a threshold $\omega_{\text{thr}} > 0$ below which the detector is blind. Thus, a final state $|Xe\gamma\rangle$ where the photon's energy is below the threshold — $\omega_\gamma < \omega_{\text{thr}}$ — will be seen by the detector as simply $|Xe\rangle$ since the photon would not be detected. In other words, *observationally* we should not classify the final states by the net number of photons regardless of their energies, however low it might be. Instead, we should count the number of *detectable* photons with $\omega_\gamma > \omega_{\text{thr}}$. As to the soft photons with $\omega_\gamma < \omega_{\text{thr}}$, the detector would not tell us if they are there or not, so **the physically measurable cross-sections should include all possible numbers of undetectably low-energy photons.** In particular,

$$\begin{aligned} d\sigma(X + e \rightarrow \text{observed } X + e) &= d\sigma(X + e \rightarrow X + e) + d\sigma(X + e \rightarrow X + e + \gamma(\omega < \omega_{\text{thr}})) \\ &\quad + d\sigma(X + e \rightarrow X + e + \gamma(\omega < \omega_{\text{thr}}) + \gamma(\omega < \omega_{\text{thr}})) + \dots, \\ d\sigma(X + e \rightarrow \text{observed } X + e + \gamma) &= d\sigma(X + e \rightarrow X + e + \gamma(\omega > \omega_{\text{thr}})) \\ &\quad + d\sigma(X + e \rightarrow X + e + \gamma(\omega > \omega_{\text{thr}}) + \gamma(\omega < \omega_{\text{thr}})) + \dots, \\ \text{etc., etc.} \end{aligned} \quad (124)$$

To the order $O(\alpha \times \sigma^{\text{tree}}(Xe \rightarrow Xe)) = O(\alpha^3)$, we should stop at one final-state photon, detectable or not, and calculate the $d\sigma(Xe \rightarrow Xe)$ to the one-loop level while the $d\sigma(Xe \rightarrow Xe\gamma)$

just to the tree level. Thus,

$$\begin{aligned}
\frac{d\sigma(X + e \rightarrow \text{observed } X + e + \gamma)}{d\Omega_e} &\approx \frac{d\sigma^{\text{tree}}(X + e \rightarrow X + e + \gamma(\omega > \omega_{\text{thr}}))}{d\Omega_e} \\
&= \int_{\omega_{\text{thr}}}^{\omega_{\text{max}}=O(E_e)} d\omega_\gamma \frac{d\sigma^{\text{tree}}(X + e \rightarrow X + e + \gamma)}{d\Omega_e d\omega_\gamma} \\
&= \frac{d\sigma^{\text{tree}}(X + e \rightarrow X + e)}{d\Omega_e} \times \frac{\alpha}{\pi} \left(2f_{\text{IR}}(q^2) \log \frac{E_e}{\omega_{\text{thr}}} + \text{finite} \right), \tag{125}
\end{aligned}$$

where the last equality comes from

$$\frac{d\sigma^{\text{tree}}(eX \rightarrow eX\gamma)}{d\Omega_e d\omega_\gamma} = \frac{d\sigma^{\text{tree}}(eX \rightarrow eX)}{d\Omega} \times \frac{\alpha}{\pi} \times \frac{2f_{\text{IR}}(q^2)}{\omega_\gamma} \quad \langle\langle \text{for } \omega_\gamma \ll E_e \rangle\rangle \tag{119}$$

and

$$\int_{\omega_{\text{thr}}}^{\omega_{\text{max}}=O(E_e)} \frac{d\omega}{\omega} = \log \frac{E_e}{\omega_{\text{thr}}} + \text{finite}. \tag{126}$$

Note that the *observed* cross-section (125) is infrared finite and does not depend on the m_γ (as long as $m_\gamma \ll \omega_{\text{thr}}$). Instead, it depends on the photon detector's low-energy threshold ω_{thr} .

Similarly, to the same order $O(\alpha \times \sigma^{\text{tree}}(Xe \rightarrow Xe)) = O(\alpha^3)$,

$$\begin{aligned}
\frac{d\sigma(X + e \rightarrow \text{observed } X + e)}{d\Omega_e} &\approx \\
&\approx \frac{d\sigma^{\text{tree}+1 \text{ loop}}(X + e \rightarrow X + e)}{d\Omega_e} + \frac{d\sigma^{\text{tree}}(X + e \rightarrow X + e + \gamma(\omega < \omega_{\text{thr}}))}{d\Omega_e} \\
&\approx \frac{d\sigma^{\text{tree}}(X + e \rightarrow X + e)}{d\Omega_e} \times \left(1 - \frac{\alpha}{\pi} \times 2f_{\text{IR}}(q^2) \log \frac{E_e}{\cancel{m_\gamma}} + \frac{\alpha}{\pi} \times \text{finite} \right) \\
&\quad + \frac{d\sigma^{\text{tree}}(X + e \rightarrow X + e)}{d\Omega_e} \times \left(0 + \frac{\alpha}{\pi} \times 2f_{\text{IR}}(q^2) \log \frac{\omega_{\text{thr}}}{\cancel{m_\gamma}} + \frac{\alpha}{\pi} \times \text{finite} \right) \\
&= \frac{d\sigma^{\text{tree}}(X + e \rightarrow X + e)}{d\Omega_e} \times \left(1 - \frac{\alpha}{\pi} \times 2f_{\text{IR}}(q^2) \log \frac{E_e}{\omega_{\text{thr}}} + \frac{\alpha}{\pi} \times \text{finite} \right). \tag{127}
\end{aligned}$$

Note how the IR regulator m_γ cancels between the two contributions to the net *observed* cross-section. Again, the observed cross-section is IR-finite, but it depends on the photon detector's threshold ω_{thr} .

Similar cancellations of IR divergences from the observed cross-sections happen at the higher loop orders. In general, to get a finite cross-section to the order $O(\alpha^L \times \sigma^{\text{tree}})$ we should combine an L -loop cross-section with no soft photons, an $L - 1$ loop cross-section with one soft photon, an $L - 2$ cross-section with 2 soft photons, *etc.*, *etc.*, ending with a tree-level cross-section with L soft photons. Individually, all these formal cross-sections are infrared-divergent, but once we combine them together into a complete observed cross-section, the IR divergence should cancel out. The *Peskin & Schroeder* textbook explains how this works in some detail in section §6.5; please read it — it's a part of your [next homework#18](#).

BREMSSTRAHLUNG

The term *bremsstrahlung* — a German word meaning *braking radiation* — originally referred to the X-rays produced by a suddenly decelerating electrons when they strike the anode. But nowadays bremsstrahlung refers to the electromagnetic radiation emitted by electrons (or perhaps other kinds of charged particles) when they suddenly change their velocities for any reason — a collision with another particle, flying into a strong magnetic field, whatever. In this section, I start with a brief classical analysis of the EM radiation by an electron abruptly changing its velocity from \mathbf{v} to \mathbf{v}' due to instantaneous scattering off some other particle X , then interpret this classical radiation in quantum terms, and eventually calculate cross-section for the soft-photon bremsstrahlung to the lowest order in QED.

Let's start with the classical analysis; I am going to be rather brief, but you can a more detailed analysis in §6.1 of the *Peskin & Schroeder* textbook. Consider an electron moving at a constant initial velocity \mathbf{v} which suddenly — in zero time — changes to \mathbf{v}' and then remains constant. In terms of the electron's trajectory $x_e^\mu(\tau)$ parametrized by the proper time τ ,

$$x_e^\mu(\tau) = \begin{cases} \tau u'^\mu & \text{for } \tau > 0, \\ \tau u^\mu & \text{for } \tau < 0. \end{cases} \quad (128)$$

The electric charge/current density due to this moving electron is

$$J^\mu(x) = (-e) \int_{-\infty}^{+\infty} d\tau \frac{dx_e^\mu}{d\tau} \times \delta^{(4)}(x_e(\tau) - x), \quad (129)$$

which Fourier transforms to

$$J^\mu(k) \stackrel{\text{def}}{=} \int d^4x e^{ikx} J^\mu(x) = (-e) \int_{-\infty}^{+\infty} d\tau \frac{dx_e^\mu}{d\tau} \times e^{ikx}. \quad (130)$$

For the electron at hand,

$$J^\mu(k) = (-e) \int_{-\infty}^0 d\tau u^\mu \times \exp(i(ku)\tau) + (-e) \int_0^{+\infty} d\tau u'^\mu \times \exp(i(ku')\tau) = ie \left(\frac{u^\mu}{(ku)} - \frac{u'^\mu}{(ku')} \right); \quad (131)$$

note conservation of this current, $k_\mu J^\mu(k) = 0$.

The EM fields emitted by this current have two components: (1) the Lorentz-boosted Coulomb fields of the initial or final electron (depending on whether the point of observation has $x^0 < +|\mathbf{x}|$ or $x^0 > +|\mathbf{x}|$); and (2) the EM radiation due to the sudden change of the electron's velocity. We are interested in the radiated field, which in momentum space is simply

$$A_{\text{rad}}^\mu(k) = \frac{1}{i|\mathbf{k}|} J^\mu(k) \times (2\pi)\delta(k^0 - |\mathbf{k}|). \quad (132)$$

Or in terms of specific radiation modes in the coordinate space,

$$A_{\text{rad}}^\mu(x) = \int \frac{d^3\mathbf{k}}{(2\pi)^3} \sum_{\lambda=\pm 1} A_{\mathbf{k},\lambda} \mathcal{E}^\mu(\mathbf{k}, \lambda) e^{-ikx} \Big|_{k^0=|\mathbf{k}|} + \text{complex conjugate} \quad (133)$$

where $\mathcal{E}^\mu(\mathbf{k}, \lambda)$ are the polarization vectors and

$$A_{\mathbf{k},\lambda} = \frac{1}{i|\mathbf{k}|} \mathcal{E}_\mu^*(\mathbf{k}, \lambda) J^\mu(\mathbf{k}, k^0 = |\mathbf{k}|). \quad (134)$$

Given the potential (133) of the radiated EM field, it's a straightforward exercise to calculate the tension fields $F_{\text{rad}}^{\mu\nu}(x)$ and the net energy they carry away. Let me skip this exercise and simply state the answer:

$$\text{EM Energy} = \int \frac{d^3\mathbf{k}}{(2\pi)^3} \sum_{\lambda} \frac{\mathbf{k}^2}{2} \times |A_{\mathbf{k},\lambda}|^2. \quad (135)$$

Moreover, thanks to $k_\mu J^\mu(k) = 0$, we have

$$\sum_{\lambda} \frac{\mathbf{k}^2}{2} \times |A_{\mathbf{k},\lambda}|^2 = \frac{1}{2} \sum_{\lambda} |\mathcal{E}_\mu^*(\mathbf{k}, \lambda) J^\mu(\mathbf{k})|^2 = -\frac{1}{2} J^\mu(\mathbf{k}) J_\mu^*(\mathbf{k}) \quad (136)$$

where the second equality is exactly similar to the way we sum QED cross-sections over the

photon's polarizations. The current $J^\mu(\mathbf{k})$ in eq. (136) is $J^\mu(k)$ for the on-shell $k^\mu = (|\mathbf{k}|, \mathbf{k})$. In terms of the wave's frequency $\omega = k^0 = |\mathbf{k}|$ and direction \mathbf{n} , eq. (131) for this current becomes

$$J^\mu(\mathbf{k}) = J^\mu(\omega, \mathbf{n}) = \frac{ie}{\omega} \left(\frac{u^\mu}{(un)} - \frac{u'^\mu}{(u'n)} \right) \quad (137)$$

where $n^\mu = (1, \mathbf{n})$ is a unit light-like vector in the direction of the wave. Consequently,

$$J^\mu(\mathbf{k})J_\mu^*(\mathbf{k}) = \frac{e^2}{\omega^2} \times \left(\frac{u^\mu}{(un)} - \frac{u'^\mu}{(u'n)} \right)^2 \quad (138)$$

where the second factor depends on the direction in which the EM radiation is emitted but not on its frequency. Altogether, this gives us

$$\begin{aligned} \text{EM Energy} &= \int \frac{d^3\mathbf{k}}{(2\pi)^3} \frac{e^2}{2\omega_k^2} \times \left[- \left(\frac{u^\mu}{(un)} - \frac{u'^\mu}{(u'n)} \right)^2 \right] \\ &= \frac{e^2}{16\pi^3} \int_0^\infty d\omega_k \omega_k^2 \int d^2\Omega(\mathbf{n}) \frac{1}{\omega_k^2} \times \left[- \left(\frac{u^\mu}{(un)} - \frac{u'^\mu}{(u'n)} \right)^2 \right]. \end{aligned} \quad (139)$$

In other words, the distribution of the bremsstrahlung energy over its frequency and direction is

$$\frac{d(\text{EM Energy})}{d\omega d\Omega(\mathbf{n})} = \frac{e^2}{16\pi^3} \times \left[- \left(\frac{u^\mu}{(un)} - \frac{u'^\mu}{(u'n)} \right)^2 \right], \quad (140)$$

and the frequency dependence cancels out from the RHS of this formula. Thus, **bremsstrahlung due to instantaneous velocity change of an electron has a completely flat frequency spectrum**. In real life, the velocity change takes a small but finite time Δt , and that cuts down on the high-frequency radiation with $\omega \gtrsim 1/\Delta t$. But in the $\delta t \rightarrow 0$ limit we consider here, the bremsstrahlung frequency spectrum is completely flat.

On the other hand, the bremsstrahlung has a strong direction dependence, as evidenced by the second factor in eq. (140). This direction dependence is particularly strong for the ultra-relativistic electrons, in which case most energy is emitted in the directions of the the initial velocity \mathbf{v} or the final velocity \mathbf{v}' . But alas, in this class I do not have time to address this issue, so let me simply

integrate the bremsstrahlung energy over the all 4π directions in which it is emitted, thus

$$\frac{d(\text{EM Energy})}{d\omega} = \frac{e^2}{4\pi^2} \times I(u, u') \quad (141)$$

where

$$I(u, u') \stackrel{\text{def}}{=} \int \frac{d\Omega(\mathbf{n})}{4\pi} \left[- \left(\frac{u^\mu}{(un)} - \frac{u'^\mu}{(u'n)} \right)^2 \right]. \quad (142)$$

Later in these notes, I shall evaluate this function of the 2 velocities and compare it to the $f_{\text{IR}}(q^2)$ which governs the IR divergence of the electrons electric form factor.

★ ★ ★

Now consider the quantum-mechanical interpretation of the flat energy spectrum (141). In QM, the EM radiation comprises a bunch of photons emitted by the suddenly accelerated electron, and their specific frequencies are subject to a probabilistic distribution. But on the average — that is, as a quantum expectation values — the numbers of photons n_γ of any particular frequencies obtain from the energy spectrum (141) as

$$\frac{d \langle n_\gamma \rangle}{d\omega_\gamma} = \frac{1}{\omega} \frac{dE_{\text{EM}}}{d\omega} = \frac{\alpha}{\pi} \times \frac{I(u, u')}{\omega_\gamma}. \quad (143)$$

Integrating over the photon's energies, we get the net number of photons we expect to emit in a bremsstrahlung event,

$$\langle N_\gamma \rangle = \frac{\alpha}{\pi} I(u, u') \times \int_{\omega_{\min}}^{\omega_{\max}} \frac{d\omega}{\omega}. \quad (144)$$

Classically, the maximal frequency here is infinite, but in the quantum theory we cannot emit a photon more energetic than the initial electron's energy. This imposes an upper limit $\omega_{\max} \lesssim E_e$ on the photon's frequency, so we do not have to worry about the UV divergence of the integral (144). On the other hand, for massless photons there is no lower limit to their frequencies, thus $\omega_{\min} = 0$ and the net expected number of photons diverges at the infrared end of the integral (144).

By itself, an infinite next expected number of very-low-energy photons is not a problem. But it becomes a problem when we try to calculate an amplitude and a cross-section for exclusive

processes involving specific numbers of emitted photons. In general, the probabilities — and hence the cross-sections — for emitting a given number n of photons are proportional to

$$\sigma(e + X \rightarrow e + X + n\gamma) \propto \langle N_\gamma \rangle^n, \quad (145)$$

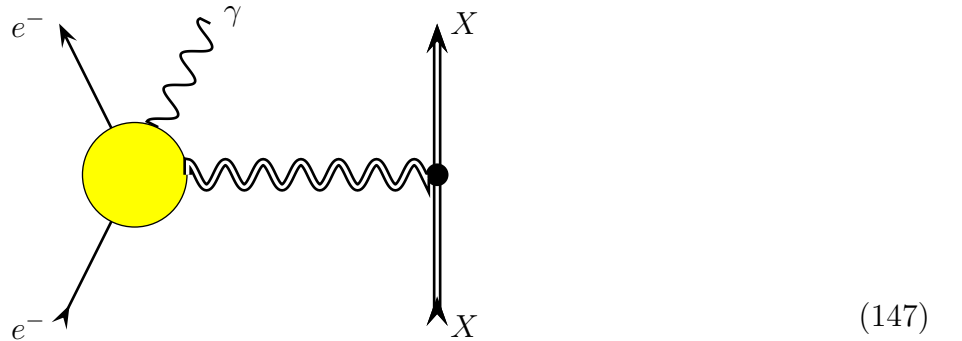
so they blow up for $\langle N_\gamma \rangle \rightarrow \infty$. Indeed, we shall see in a moment that for a hard Coulomb scattering of an electron off a heavy pointlike target X , the probability of emitting exactly one extra photon is

$$\sigma(e + X \rightarrow e + X + \gamma) \propto \langle N_\gamma \rangle \sim \alpha \times \int_0^{\omega_{\max}} \frac{d\omega}{\omega}, \quad (146)$$

which suffers from the IR divergence as $\omega \rightarrow 0$.

* * *

With all of the above in mind, let's go back to QED and calculate the tree-level cross-section for the inelastic electron scattering which emits one photon:



In particular, let's verify the relation

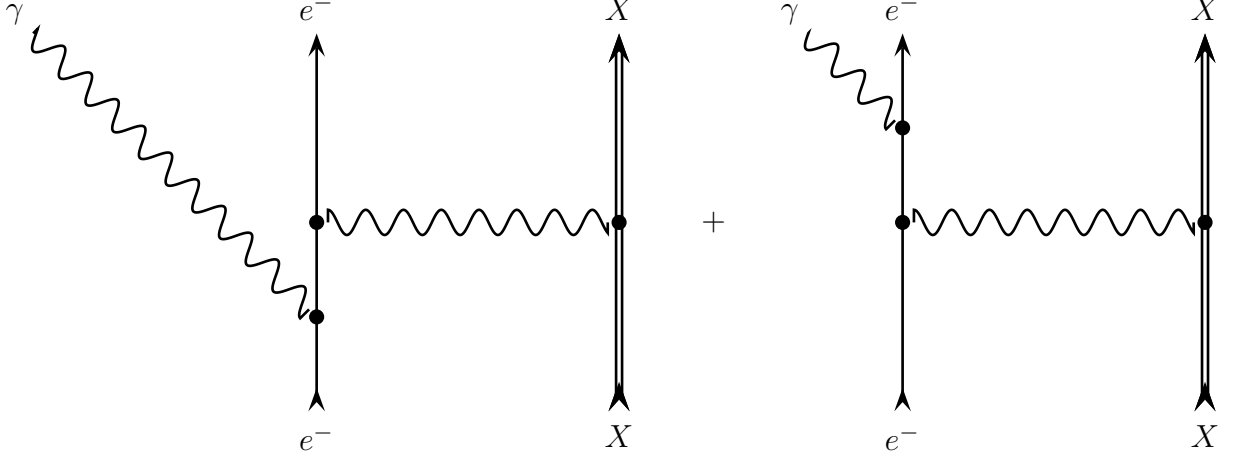
$$\frac{d\sigma^{\text{tree}}(eX \rightarrow eX\gamma)}{d\Omega_e d\omega_\gamma} = \frac{d\sigma^{\text{tree}}(eX \rightarrow eX)}{d\Omega} \times \frac{\alpha}{\pi} \times \frac{2f_{\text{IR}}(q^2)}{\omega_\gamma} \quad (119)$$

where f_{IR} is exactly as in eq. (102) for the coefficient of the IR divergence on the one-loop electric form factor. Along the way, we shall also relate this $f_{\text{IR}}(q^2)$ to the classical factor $I(u, u')$ of the

bremsstrahlung intensity, specifically

$$I(u = p/m_e, u' = p'/m_e) = 2f_{\text{IR}}(q^2). \quad (148)$$

At the tree level, there are 2 diagrams contributing to the inelastic scattering (147), namely



Evaluating them for a nearly-recoil-less X particle, we get the amplitude

$$\frac{\mathcal{M}_{\text{tree}}}{2M_X} = \frac{e^2}{q^2} \times \left(\bar{u}(p')\gamma^0 \times \frac{i}{\not{p} - \not{k} - m} (ie \not{\epsilon}^*) u(p) + \bar{u}(p')(ie \not{\epsilon}^*) \frac{i}{\not{p}' + \not{k} - m} \times \gamma^0 u(p) \right) \quad (149)$$

where ϵ^μ is the polarization vector of the outgoing photon. For the on-shell incoming electron and outgoing photons

$$(p - k)^2 - m^2 = k^2 - 2(pk) + p^2 - m^2 = -(2pk), \quad (150)$$

hence

$$\begin{aligned} \frac{i}{\not{p} - \not{k} - m} (ie \not{\epsilon}^*) u(p) &= e \frac{(\not{p} - \not{k} + m) \not{\epsilon}^*}{2(pk)} u(p) \\ &= e \frac{2(p\epsilon^*) + \not{\epsilon}^*(m - \not{p}) - \not{k} \not{\epsilon}^*}{2(pk)} u(p) \\ &= e \frac{2(p\epsilon^*) - \not{k} \not{\epsilon}^*}{2(pk)} u(p) \quad \langle\langle \text{since } (\not{p} - m)u(p) = 0 \rangle\rangle \\ &\approx e \frac{2(p\epsilon^*)}{2(pk)} u(p) \quad \langle\langle \text{because } k \ll p \rangle\rangle. \end{aligned} \quad (151)$$

and therefore

$$\bar{u}(p')\gamma^0 \times \frac{i}{\not{p}' - \not{k} - m}(ie\not{\epsilon}^*)u(p) \approx \bar{u}(p')\gamma^0 u(p) \times e\frac{(p\epsilon^*)}{(pk)}. \quad (152)$$

In a similar way, for the second tree diagram we get

$$(\not{p}' + k)^2 = +2(p'k), \quad (153)$$

hence

$$\begin{aligned} \bar{u}(p')(ie\not{\epsilon}^*)\frac{i}{\not{p}' + \not{k} - m} &= -e\bar{u}(p')\frac{\not{\epsilon}(\not{p}' + \not{k} + m)}{2(p'k)} \\ &= -e\bar{u}(p')\frac{2(p'\epsilon^*) + (m - p')\not{\epsilon}^* + \not{\epsilon}\not{k}}{2(p'k)} \\ &= -e\bar{u}(p')\frac{2(p'\epsilon^*) + \not{\epsilon}\not{k}}{2(p'k)} \quad \langle\langle \text{since } \bar{u}(p') \times (m - p') = 0 \rangle\rangle \\ &\approx -e\bar{u}(p') \times \frac{(p'\epsilon^*)}{(p'k)} \quad \langle\langle \text{since } k \ll p' \rangle\rangle, \end{aligned} \quad (154)$$

and therefore

$$\bar{u}(p')(ie\not{\epsilon}^*)\frac{i}{\not{p}' + \not{k} - m} \times \gamma^0 u(p) \approx \bar{u}(p')\gamma^0 u(p) \times \frac{-(p'\epsilon^*)}{(p'k)}. \quad (155)$$

Altogether, the complete tree amplitude due to both diagrams is

$$\begin{aligned} \mathcal{M}_{\text{tree}}(e^- X \rightarrow e^- X \gamma) &\approx \mathcal{M}_{\text{tree}}(e^- X \rightarrow e^- X) \times e \left(\frac{(p\epsilon^*)}{(pk)} - \frac{(p'\epsilon^*)}{(p'k)} \right) \\ &= \mathcal{M}_{\text{tree}}(e^- X \rightarrow e^- X) \times \frac{e}{\omega} \left(\epsilon_\mu^* \left(\frac{p'^\mu}{(np')} - \frac{p^\mu}{(np)} \right) \right) \end{aligned} \quad (156)$$

where $\omega = k^0$ is the photon's frequency and n^μ is the unit 4-vector for its direction, $n^0 = |\mathbf{n}| = 1$.

To convert the amplitude (156) into a cross-section, we need to sum $|\mathcal{M}|^2$ over photon polarizations, sum/average over the electron polarizations, and take care of the phase space factors. Fortunately, the sum/average over the electron spins works in exactly the same way with or with-

out the photon, thus

$$\begin{aligned}
\overline{|\mathcal{M}|^2}_{\text{tree}}(e^- X \rightarrow e^- X \gamma) &= \overline{|\mathcal{M}|^2}_{\text{tree}}(e^- X \rightarrow e^- X) \times \frac{e^2}{\omega^2} \sum_{\lambda} \left| \epsilon_{\mu}^*(k, \lambda) \left(\frac{p'^{\mu}}{(np')} - \frac{p^{\mu}}{(np)} \right) \right|^2 \\
&= \overline{|\mathcal{M}|^2}_{\text{tree}}(e^- X \rightarrow e^- X) \times \frac{e^2}{\omega^2} \left[- \left(\frac{p'}{(np')} - \frac{p}{(np)} \right)^2 \right].
\end{aligned} \tag{157}$$

Furthermore, for $\omega \ll E_e$ the electron's phase-space factors also do not care about the extra soft photon, hence

$$d\sigma_{\text{tree}}(e^- X \rightarrow e^- X \gamma) = d\sigma_{\text{tree}}(e^- X \rightarrow e^- X) \times \frac{\omega^2 d\omega d\Omega_n}{(2\pi)^3 2\omega} \times \frac{e^2}{\omega^2} \left[- \left(\frac{p'}{(np')} - \frac{p}{(np)} \right)^2 \right]. \tag{158}$$

Note that the last factor inside the $[\dots]$ here is exactly as in the classical bremsstrahlung spectrum (140), since for $p = m_e u$ and $p' = m_e u'$

$$\frac{p^{\mu}}{(pn)} = \frac{u^{\mu}}{(un)} \quad \text{and} \quad \frac{p'^{\mu}}{(p'n)} = \frac{u'^{\mu}}{(u'n)}. \tag{159}$$

Consequently, when we integrate the partial cross-section (158) over the photon's directions — but not their energies or the electron's directions, — we get

$$\frac{d\sigma^{\text{tree}}(eX \rightarrow eX \gamma)}{d\Omega_e d\omega_{\gamma}} = \frac{d\sigma^{\text{tree}}(eX \rightarrow eX)}{d\Omega} \times \left(\frac{e^2}{4\pi^2} = \frac{\alpha}{\pi} \right) \times \frac{I(p', p)}{\omega_{\gamma}} \tag{160}$$

where

$$I(p, p') = I(u, u') = \int \frac{d\Omega(n)}{4\pi} \left[- \left(\frac{p'}{(np')} - \frac{p}{(np)} \right)^2 \right] \tag{161}$$

exactly as in the classical formula (142). In a moment we shall see that $I(p, p') = 2f_{\text{IR}}(q^2)$, and that turns eq. (160) into eq. (119) with the precisely right coefficient of the IR divergence to cancel the IR divergence of the one-loop photon-less cross-section.

So let us evaluate the angular integral (161). Our first step in that direction is to expand the Lorentz square in the integrand:

$$-\left(\frac{p'}{(np')} - \frac{p}{(np)}\right)^2 = -\frac{(p^2 = m^2)}{(np)^2} - \frac{(p'^2 = m^2)}{(np')^2} + \frac{2(pp') = 2m^2 - q^2}{(np)(np')}. \quad (162)$$

Next, we integrate each of these three terms over the directions of the unit 3-vector \mathbf{n} . For the first term, we have

$$\begin{aligned} \int \frac{d^2\Omega_{\mathbf{n}}}{4\pi} \frac{1}{(np)^2} &= \frac{2\pi}{4\pi} \int_{-1}^{+1} d\cos\theta \frac{1}{(E - |\mathbf{p}| \times \cos\theta)^2} = \frac{1}{2} \times \frac{1}{|\mathbf{p}|} \int_{E-|\mathbf{p}|}^{E+|\mathbf{p}|} \frac{d(E - |\mathbf{p}| \times \cos\theta)}{(E - |\mathbf{p}| \times \cos\theta)^2} \\ &= \frac{1}{2|\mathbf{p}|} \left(\frac{1}{E - |\mathbf{p}|} - \frac{1}{E + |\mathbf{p}|} \right) = \frac{1}{E^2 - |\mathbf{p}|^2} = \frac{1}{m^2}. \end{aligned} \quad (163)$$

Likewise, for the second term

$$\int \frac{d^2\Omega_{\mathbf{n}}}{4\pi} \frac{1}{(np')^2} = \frac{1}{m^2}. \quad (164)$$

The third term is trickier, so let's use the Feynman parameter trick:

$$\frac{1}{(np) \times (np')} = \int_0^1 d\xi \frac{1}{[(1-\xi) \times (np) + \xi \times (np')]^2} = \int_0^1 d\xi \frac{1}{(np_\xi)^2} \quad (165)$$

where

$$p_\xi^\mu = (1-\xi) \times p^\mu + \xi \times p'^\mu. \quad (166)$$

Consequently,

$$\int \frac{d^2\Omega_{\mathbf{n}}}{4\pi} \frac{1}{(np) \times (np')} = \int_0^1 d\xi \int \frac{d^2\Omega_{\mathbf{n}}}{4\pi} \frac{1}{(np_\xi)^2} \quad (167)$$

where

$$\int \frac{d^2\Omega_{\mathbf{n}}}{4\pi} \frac{1}{(np_\xi)^2} = \frac{1}{p_\xi^2} \quad (168)$$

exactly as in eqs. (163) or (164), except for $p_\xi^2 \neq m^2$ but rather

$$\begin{aligned}
p_\xi^2 &= [(1-\xi)p + \xi p']^2 \\
&= (1-\xi)^2 \times (p^2 = m^2) + \xi^2 \times (p'^2 = m^2) + \xi(1-\xi) \times (2pp' = 2m^2 - q^2) \\
&= m^2 - \xi(1-\xi)q^2.
\end{aligned} \tag{169}$$

Altogether,

$$\int \frac{d^2\Omega_{\mathbf{n}}}{4\pi} \frac{1}{(np) \times (np')} = \int_0^1 \frac{d\xi}{m^2 - \xi(1-\xi)q^2}. \tag{170}$$

Finally, plugging eqs. (163), (164), and (170) back into eqs. (142) and (162), we arrive at

$$I(p', p) = -\frac{m^2}{m^2} - \frac{m^2}{m^2} + (2m^2 - q^2) \times \int_0^1 \frac{d\xi}{m^2 - \xi(1-\xi)q^2} = \int_0^1 d\xi \left(\frac{2m^2 - q^2}{m^2 - \xi(1-\xi)q^2} - 2 \right). \tag{171}$$

Finally comparing the last formula with the definition of the f_{IR} ,

$$f_{\text{IR}}(q^2/m_e^2) = \int_0^1 d\xi \left(\frac{2m_e^2 - q^2}{2[m^2 - \xi(1-\xi)q^2]} - 1 \right), \tag{102}$$

we immediately see that $I(p', p) = 2f_{\text{IR}}(q^2/m_e^2)$, *quod erat demonstrandum*.