Quantum Mechanics and the Warping of Spacetime: Infinite Superpositions and the Schwarzchild Metric.

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Abstract

This paper proposes new physics. In a radically different approach, it forms fundamental particles as infinite superpositions, but all with non zero mass borrowed from a Higgs type scalar field. However energy is also borrowed from zero point vector fields. Just as the Standard Model divides the fundamental particles into two types…those with mass and those without with the Higgs mechanism providing the difference…infinite superpositions seem also to divide naturally into two sets: (a) those with "infinitesimal" mass, and (b) those with significant mass (from micro electron volts upwards). In the infinitesimal set (a), photons, gluons and gravitons (to fit with cosmology and the expansion of the cosmos) all have $\approx 10^{-34}$ eV mass, approximately the inverse of the causally connected horizon radius. These values are so close to zero the Standard Model symmetry breaking approach remains essentially valid. Particles travelling this close to the speed of light have virtually fixed helicity with the Higgs mechanism increasing their mass from infinitesimal type (a) to significant or measureable type (b) values. Also the energy in the zero point fields (borrowed to build the fundamental particles) is limited, particularly at the extreme wavelengths of virtual gravitons interacting at near horizon radii. Any causally connected region grows with time after the big bang and the number of virtual gravitons with wavelengths similar to the size of the causally connected region increases approximately as the square of the causally connected mass. Space has to expand progressively with time after the big bang, increasing the zero point energy available to meet this increased requirement. For similar reasons the extra gravitons near mass concentrations change the metric in proportion to m/r, in accordance with the Schwarzschild solution of Einstein's equations. It suggests possible spacetime boundaries at the event horizons of black holes in line with one of the current Firewall Paradox implications. Approximately the first two thirds of this paper build and analyse the fundamental particles from infinite virtual superpositions. The final third looks at the expanding Universe and connections with General Relativity, but only after attempting to show that infinite superpositions are equivalent to the Standard Model fundamental particles, apart from infinitesimal differences.

1 Introduction

Since the weak and electromagnetic forces were unified in the 1960's physicists have wished to somehow unite all the fundamental forces. Initially it seemed that three of the forces (excluding gravity) could unite near the Planck scale. High energy experiments however ruled out the possibility of single energy unification. Supersymmetry was proposed as a possible solution addressing several problems, but also modifying high energy running constants of the three forces in such a manner they united near the Planck scale. None of the particles predicted however have yet been seen. String theory is currently seen by many as the most likely future path, but not all physicists are comfortable with its non-testability and need for 10 or 11 dimensions. The enormous landscape of different universes or multiverse it proposes is also widely regarded as the solution to the minute amount of dark energy proposed to explain the current accelerating expansion of the universe. Recent cosmological surveys [1] however have not so far supported the dark energy explanation. This all suggests some important and relevant questions, for example:

1. Is it possible that the fundamental forces may connect in some different way?
2. Are the extra dimensions of Supersymmetry and String Theory really necessary?
3. Is "The Multiverse" the only explanation of accelerating cosmic expansion?
4. Can the problems these theories were meant to solve be addressed differently?

Approaching all this in a completely new direction, this paper attempts to address these questions using basic principles of quantum mechanics and relativity. Apart from infinitesimal differences it seems to be consistent with the Standard Model. It requires the universe to expand after the big bang and possibly in an accelerating manner. It changes the metric around mass concentrations in accordance with General Relativity. It requires photons, gluons and gravitons to have a mass of $\approx 10^{-34} eV$, very close to some other recent proposals [2] [3] giving gravitons a mass of $< 10^{-33} eV$ to explain the accelerating expansion of the universe. It argues that General Relativity cuts off at the Black Hole event horizon, one of the possible implications of the current Firewall paradox.[14] [15] [16] [17] [18].

1.1 Summary

Papers modifying the Standard Model are too numerous to list, however we briefly touch on a small number of some early versions of these in section 1.1.2. The approach in this paper is very different from that in most of these earlier papers, with the main differences summarized below.

1.1.1 General Relativity as our starting point

General Relativity tells us that all forms of mass, energy and pressure are sources of the gravitational field. Thus to create gravitational fields all spin ½ leptons & quarks, spin 1 gluons, photons, $W^{\pm}$ & Z^0 particles etc. emit virtual gravitons.

The starting point of this paper assumes there is a common thread uniting these fundamental particles making this possible. Equations are developed that unite the amplitudes of the colour and electromagnetic coupling constants with that of gravity. The precision required by quantum mechanics for half integral and integral angular momentum allows gravity to be included, despite the vast disparity in magnitude between gravity and the other two. This combination of colour, electromagnetic and gravitational amplitudes in the same equation is possible only because of a radically different approach taken in this paper: An approach using infinite superpositions of positive and negative integral $\hbar$ angular momentum virtual wavefunctions for spin ½, spin 1 and spin 2 particles. The final result is almost identical to the Standard Model, with infinitesimal but important differences.

The total angular momentum can be summed over all wavenumbers k; from $k = 0$ to some cutoff value k_{cutoff}. We will assume (as with many unification theories) that the cutoff for these infinite superpositions is somewhere near Planck scale. Firstly imagine a universe where the gravitational constant $G \to 0$. As $G \to 0$, the Planck length $L_P \to 0$, the Planck energy $E_P \to \infty$ and $k_{cutoff} \to \infty$ also. If we sum the angular momentum of these infinite superpositions when $G \to 0$ (i.e. from $k = 0$ to $k_{cutoff} \to \infty$) we get precisely half integral or integral $\hbar$ for the fundamental spin ½, spin 1 & spin 2 particles in appropriate m states. If we now put $G > 0$ the infinitesimal effect of including gravity can be balanced by an equal but opposite effect due to the non-infinite cutoff value in k. A near Planck scale superposition cutoff requires gravity to be included to get precisely half integral or integral $\hbar$. (Section 4.2)

These infinite superpositions have another very relevant property relating to the fact that all experiments indicate that fundamental particles such as electrons behave as point particles. Each wavefunction with wavenumber k, which we label as ψ_k, has a maximum radial probability at $r \approx 1/k$ and they all look the same (Figure 1.1. 1.) Every wavefunction ψ_k of these infinite superpositions, interacts only with virtual photons (for example) of the same k; if superpositions representing say an electron are probed with such photons (that interact only with wavefunction ψ_k) the resolution possible is of the same order as the dimensions of ψ_k, both have $r \approx 1/k$. The higher the energy of the probing particle the smaller the ψ_k it interacts with, the resolution of an observing photon can never be fine enough to see any ψ_k

dimensions. Even if this energy approaches the Planck value, with a matching ψ_k radius near the Planck length it is still not possible to resolve it. This behaviour is consistent with the quantum mechanical properties of point particles.

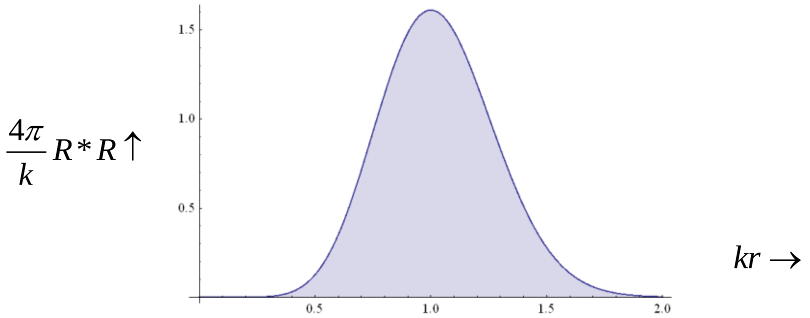


Figure 1.1. 1 The radial probability of the dominant $n=6$ for spin ½ wavefunction ψ_{6k}.

1.1.2 Primary interactions and Secondary interactions

Supposing that superpositions can in fact build the fundamental spin ½, spin 1, and spin 2 particles, then what builds the superpositions? Before answering that question, in this paper we divide the world of all interactions into two categories.

Secondary Interactions are those we are familiar with and are covered by the Standard Model, but with the addition of gravity, which is not included in the Standard Model. They take place between the fundamental spin ½, spin 1 and spin 2 particles formed from infinite superpositions. They are the QED/QCD etc. interactions of all real world experiments.

Primary Interactions on the other hand are those that build infinite superpositions and are hidden to the real world of experiments.

The majority of this paper is about these primary interactions, and the superpositions they build representing the fundamental spin ½, spin 1 and spin 2 particles. Primary interactions are between spin zero particles borrowed from a Higgs type scalar field and the zero point vector fields. In the 1970's models were proposed with preons as common building blocks of leptons and quarks [4] [5] [6] [7]. In contrast with the virtual particles in this paper some of these earlier models used real spin ½ building blocks. Real substructure has difficulties with large masses if compressed into the small volumes required to approach point particle behaviour. On the other hand with *virtual substructure borrowing energy from zero point fields the mass contribution at high* k *values can be cancelled* (section 3.2.1). As in earlier models this paper also calls the common building blocks preons, *but here the preons are both virtual and spin zero.* They also now build all spin ½ leptons and quarks, spin 1 gluons, photons, W & Z particles, plus spin 2 gravitons in contrast to only the leptons and quarks in the earlier models. (See Table 2.2. 1.)

As these preons have zero spin they possess no weak charge, primary interactions (section 2.2.1) can take place only with the zero point colour, electromagnetic and gravitational fields. The three primary coupling constants for each of these three zero point fields are different from (but related to) the secondary coupling constants. The behaviour of primary coupling is also entirely different from secondary coupling. Secondary coupling strengths vary (or run) with wavenumber k (the electromagnetic increasing with k and colour decreasing with k). In contrast, primary coupling strengths (or constants) do not run. In this paper virtual preons are continually born out of a Higgs type scalar field, existing only for time $\Delta t \approx \hbar / E.$ At their birth, they interact while still bare with zero point vector fields at this instant of birth $t = 0$. The primary coupling constants consequently are fixed for all k: there is no time for charge canceling or reinforcing, which in secondary interactions forms around the bare charge progressively after its birth. The equations work only if this is true, and they also work only if the primary colour coupling constant is 1. This does not seem implausible as it simply means that primary colour coupling is certain (sections 2.2.2). The ratio between the primary and secondary colour coupling constants labelled χ_c is thus (*if primary colour coupling is 1*) the inverse of the secondary (or usual QCD $/\alpha_3^{-1}$) colour coupling constant *at the superposition cutoff near the Planck scale.* (Sections 3.3, 4.1.1 & 4.3)

To enable the primary coupling to colour, electromagnetic and gravitational zero point fields, preons need colour, electric charge and mass. Red green or blue coloured preons have positive electric charge; anticolour red, green or blue preons have negative electric charge. Their mass which is borrowed from some type of scalar Higg's field must always be non-zero, which is discussed further in section 1.1.3. As there are 8 gluon fields, superpositions are built with 8 virtual preons for each virtual wavefunction ψ_k. The nett sum of these 8 electric charges is $0, \pm 2, \pm 4, \pm 6$, and never $> \pm 6$. This leads to the usual $0, \pm 1/3, \pm 2/3, \pm 1$ electric charge seen in the real world. Various combinations of these 8 preons in appropriate superpositions can build leptons and quarks, colour changing and neutral gluons, neutral photons, neutral massive Z^0 photons and the charged massive $W^{\pm}$ photons. (Table 2.2. 1)

1.1.3 Photons, gluons and gravitons with infinitesimal mass ($\approx 10^{-34} eV$).

For many decades after the discovery of the neutrino in the 1930s it was thought to be massless, and to travel at velocity c. Towards the end of last century however evidence slowly accumulated that this may not in fact be quite true, and that the family of 3 neutrinos have masses in the electron volt range. Due to this very low mass, and their normal emitted energies, they invariably travel at virtually the velocity of light c. Photons also have always been seen as massless traveling precisely at velocity c, except in the case of the massive $W^{\pm}$ & Z^{0}. Massless virtual photons have an infinite range, which has always been seen as an absolute requirement of the electromagnetic field. On the other hand, this paper requires some rest frame (even if this frame normally moves virtually at c) in which to build all the fundamental particles. Table 6.2 1 suggests photons, gluons and gravitons have $\approx 10^{-34} eV$ mass with *a range of approximately the inverse of the causally connected horizon radius,* and velocities sufficiently close to that of light their helicity remains essentially fixed. This allows some form of Higgs mechanism to increase this infinitesimal mass to the various values in the massive set. These infinitesimal masses are in line with other recent proposals [2] [3] giving gravitons a mass of $<10^{-33} eV$ to explain accelerating expansion.

The virtual wavefunction we use is $\psi_{nk} = C_{nk} r^3 \exp(-n^2 k^2 r^2 / 18) Y(\theta, \varphi)$ an $l = 3$ wavefunction. This virtual $l = 3$ property is normally hidden. In the same way as scattering experiments on spin 0 pions show spin 0 properties, and not the properties of the two canceling spin ½ component particles, this $l = 3$ property of the virtual components of superpositions is not visible in the real world. Scattering experiments can exhibit only the spin properties of the resulting particle. The individual angular momentum vectors $|\mathbf{L}| = 2\sqrt{3}\hbar$ of the infinite superposition all sum to a resulting: $|\mathbf{L}_{Total}| = (\sqrt{3}/2)\hbar$, $\sqrt{2}\hbar$ or $\sqrt{6}\hbar$ for spin ½ , spin 1 or spin 2 respectively, in a similar way to two spin ½ particles forming spin 0 or spin 1 states.

The wavefunction $\psi_{nk} = C_{nk} r^3 \exp(-n^2 k^2 r^2 / 18) Y(\theta, \varphi)$ has Eigenvalues $\mathbf{P}_{nk}{}^2 = n^2 \hbar^2 k^2$ with $|\mathbf{P}_{nk}| = n\hbar k$, suggesting it borrows n parallel $|\hbar \mathbf{k}|$ quanta from zero point vector fields provided n is integral. We can see this by letting $k \to \infty$ allowing energy $E \to n\hbar\omega$ by absorbing n quanta $\hbar\omega$ from the zero point vector fields (section 2.3.2). As spin 3 needs at least 3 spin 1 particles to create it, the lowest integral number that n can be is 3. The virtual $l = 3$ property can however be used to derive the magnetic moment of a charged spin ½, $m = \pm 1/2$ state as a function of n. Section 3.5 shows $g = 2$ Dirac electrons need an average (over integral n states) of $\bar{n} \approx 6.0135$. Three member superpositions $\psi_k = \sum c_{nk} \psi_{nk}$ with $n = 5, 6, \&7$ achieve this, creating Dirac spin ½ states. We also find that $n = 6$ is the dominant member and each superposition ψ_k needs at least 3 members to make all the equations consistent for Dirac particles. Secondary interactions at any wavenumber k can occur with ψ_k if integers n

change by ± 1, thus changing the Eigenvalues $|\mathbf{P}| = n\hbar k$ by $\pm\hbar k$ where this can be only a temporary rearrangement of the triplets of values of n. This is true, whether the interaction is with leptons, quarks, photons, gluons, W & Z particles, or gravitons. (Section 3.3)

1.1.4 Superpositions require only squared vector potentials

The wavefunction $\psi_{nk} = C_{nk} r^3 \exp(-n^2k^2r^2/18)Y(\theta,\varphi)$ also requires a squared vector potential to create it: $Q^2A^2 = n^4\hbar^2k^4r^2/81$. There are no linear potential terms in contrast with secondary interactions. In our momentum operator for primary interactions we thus use $\hat{P}^2 = -\hbar^2\nabla^2 + Q^2A^2$, no linear potential terms are included and Q simply represents a collective symbol for all the effective charges concerned. As an example, the dominant $n=6$ wavefunction of a spin ½ Dirac ψ_k requires a squared vector potential of $Q^2A^2 = n^4\hbar^2k^4r^2/81 = 16\hbar^2k^4r^2$ (section 2.3.1). Primary coupling between the 8 virtual preons and the colour, electromagnetic and gravitational zero point fields produces a vector potential squared value for all infinite superpositions which can be expressed as:

$$Q^2A^2 = \frac{\left[8 + 8\sqrt{\alpha_{EMP}} + im_0(1+\beta_{nk}{}^2)\sqrt{G_p/(2s\hbar c)}\right]^2(\hbar^2k^4r^2)}{3\pi(sN)(1+\varepsilon)}\left[\frac{(sN)(1+\varepsilon)dk}{k}\right]$$

(Where the length of the complex vector is squared here.) The significance of the cancelling top and bottom factors (sN) is explained in section 2.1.2. Also the cancelling $(1+\varepsilon)$ factors are due to gravity and explained in section 4.2. The primary_colour coupling amplitude is 1 to each of the eight preons, and $\sqrt{\alpha_{EMP}}$ the primary electromagnetic coupling. This equation applies regardless of the individual preon colour or electric charge signs, whether positive or negative (section 2.2.3). The primary gravitational coupling is to the particle mass m_0, and a pressure term $m_0\beta_{nk}{}^2$. The primary gravitational constant is G_P divided by $\hbar c$ to put it in the same form as the other two coupling constants. The magnitude of the total angular momentum vector of the infinite superposition is $|\mathbf{L}_{Total}| = \sqrt{s(s+1)}$.) This Q^2A^2 without the gravity term generates superpositions with probability $(N\cdot s)dk/k$ where s is the superposition spin, $N=1$ for massive spin ½ superpositions but $N=2$ for both massive photon superpositions and infinitesimal mass superpositions (Table 4.3. 1, section 6 and its subsections cover this more fully). Section 4.2 includes gravity raising the superposition probability to $(1+\varepsilon)(N\cdot s)dk/k$ where the infinitesimal ε (not to be confused with infinitesimal mass) is $\varepsilon \approx m_0{}^2G_P/(110s\hbar c) \approx 10^{-45}$ for electrons, and $\varepsilon \approx 10^{-34}$ for a Z^0 The ψ_k superpositions require at least three integral n members. The following three member superpositions seem to fit the Standard Model best (see Table 4.3. 1)

Spin ½ massive $N=1$ fermion superpositions $\psi_k = \sum_{n=5,6,7} c_{nk}\psi_{nk}$.

Spin 1 massive $N=2$ boson superpositions, also spins 1 & 2 $\psi_k = \sum_{n=3,4,5} c_{nk}\psi_{nk}$.

infinitesimal mass $N=2$ superpositions

Below are infinite superpositions $\psi_{\infty,s,m}$ for only *massive spins* ½ & 1. The symbol ∞ refers to the infinite sum, s the spin of the resulting real particle, m its angular momentum state, and ss a spherically symmetric state. Section 3.1.3 explains this format. Also square cutoffs in wavenumber k are used here for simplicity; exponential cutoffs in k are introduced in section 4.3. Infinitesimal mass superpositions are introduced in section 6.2.

$$\begin{aligned}
&\text{Massive } N=1 \text{ Spin } \tfrac{1}{2}, \psi_{\infty,1/2,m} = \sum_{n=5,6,7} c_n \int_0^{k(cutoff)} \left[\frac{(\psi_{nk,ss})}{\gamma_{nk}} + \beta_{nk}(\psi_{nk,4m}) \right] \sqrt{\frac{1+\varepsilon}{2k}}dk \\
&\text{Massive } N=2 \text{ Spin } 1, \psi_{\infty,1,m} = \sum_{n=3,4,5} c_n \int_0^{k(cutoff)} \left[\frac{(\psi_{nk,ss})}{\gamma_{nk}} + \beta_{nk}(\psi_{nk,2m}) \right] \sqrt{\frac{2(1+\varepsilon)}{k}}dk
\end{aligned} \qquad (1.1.\ 1)$$

In these infinite superpositions the probability that the wavefunction is spherically symmetric is $1/\gamma_{nk}{}^2 = 1-\beta_{nk}{}^2$ and the probability that it is an m state is $\beta^2{}_{nk}$ where β_{nk} is the magnitude of the velocity of the centre of momentum of the primary interactions that generate each ψ_{nk} . This is similar to the superposition of time and spatially polarized virtual photons in QED. For example spin ½ has probabilities of $1/\gamma_{nk}{}^2 = 1-\beta_{nk}{}^2$ spherically symmetric ψ_{nk} wavefunctions, and $\beta^2{}_{nk}\ \times (\psi_{nk}, m=\pm 2)$ wavefunctions. Each ψ_k is normalized to 1 but the infinite superpositions $\psi_{\infty,s,m}$ are not normalized, diverging logarithmically with k ; the same logarithmic divergence that applies to virtual photon emission. *Real wavefunctions have to be normalized to one as they refer to finding a real particle somewhere but this need not apply here.* Each member of these spin ½ superpositions has probability $dk(1+\varepsilon)/2k$, and if electrically charged emits virtual photons with probability $4\alpha/\pi$. Ignoring the factor of $(1+\varepsilon) \approx 1+ \approx 10^{-45}$, the overall virtual scalar photon emission probability is the usual $(2\alpha/\pi)dk/k$. (Possible implications of the infinitessimal ε are discussed in section 6.6) We also find in section 3.1 that $m=+2$ virtual wavefunctions have $\beta^2{}_{nk}$ probability of leaving an $m=-2$ debt in the zero point fields. Integrating this over all k produces a total angular momentum for a spin ½ state of $(\hbar/2)(1-1/\gamma^2{}_{Cutoff})(1+\varepsilon)$, (section 3.2.2). When $1/k_{Cutoff}$ is near the Planck length $(1-1/\gamma^2{}_{Cutoff}) = 1/(1+\varepsilon)$. A similar integration over all k for the rest energy of the infinite superposition also leads to $\pm m_0 c^2 (1-1/\gamma^2{}_{Cutoff})(1+\varepsilon)$, (section 3.2.1).

The infinitesimal quantity ε vanishes in a zero gravity, zero Planck length universe where k_{Cutoff} & $\gamma_{Cutoff} \to \infty$. In this paper each preon borrowed from a Higgs type scalar field has virtual rest mass. The superposition mass/energy is obtained by summing squared momenta over all k. The equations are based on probabilities of these in a similar manner to those for angular momentum. This suggests the superposition or equivalent particle mass is both energy and mass borrowed from zero point *vector,* and Higgs type *scalar* fields.

The final sections of this paper (5 & 6) argue that the limited zero point energies available at the extreme wavelengths of virtual gravitons spanning the causally connected cosmos require it to continually expand and possibly in an accelerating manner. Sections 5.2 & 5.2.2 argue that the warping of spacetime around mass concentrations is consistent with local observers measuring a constant background density of virtual gravitons $\rho_{Gk\min}$ at the wavelength limit. This can only happen if at any radius *r* around a mass *m*, space expands proportionally to *m/r* in accordance with the Schwarzschild solution. We argue that this implies General Relativity and the warping of spacetime is a consequence of Quantum Mechanics.

The first two thirds of this paper are about the primary interactions between spin zero preons and spin one quanta that build the fundamental particles. The Standard Model is about the secondary interactions between them. *The weak force is only between spin ½ particles and thus a secondary interaction. It can not be involved in primary interactions*. Apart from infinitesimal effects, such as infinitesimal masses, the properties of fundamental particles covered in this paper seem consistent with their Standard Model counterparts. This paper relates the colour coupling constant of secondary (Standard Model) interactions at the near Planck length cutoff to the secondary (Standard Model) interaction electromagnetic coupling constant also at this cutoff. Section 4.3 suggests gravity cuts off interactions exponentially somewhere between ≈ 2 and $2.1\times10^{18}GeV$. or $\approx 1/6$ of the Planck energy. The secondary electromagnetic coupling constant suggested at this cutoff is $\alpha^{-1}{}_{EMS} = \alpha^{-1}{}_{EM} \approx 106$, approximately in line with Standard Model predictions assuming three families of fermions and one Higgs field. (See Figure 4.1. 1 & Figure 4.1. 2). At this cutoff the Standard Model secondary colour coupling constant is $\alpha^{-1}{}_{3} \approx 50.4$ which is also the primary to secondary coupling ratio χ. Only after attempting to show that infinite superpositions can be equivalent to the Standard Model particles do we try to connect them with General Relativity.

2 Building Infinite Virtual Superpositions

2.1 The possibility of Infinite Superpositions

2.1.1 Early ideas

After World War II there was still much confusion about QED. In 1947 at the Long Island Conference the results of the Lamb shift experiment were announced [8]. Some of the first early explanations that gave approximately correct answers used simple semi classical thinking to get a better understanding of what seemed to be going on. These early ideas helped to eventually lead to the QED of today, perhaps in a similar manner to the way Bohr's original simple semi classical explanation of quantized atomic energy levels played such a large part in the eventual development of full three dimensional wavefunction solutions of atoms, and quantum mechanics. We start this paper with an example of a semi classical Lamb shift explanation that seems to lead into the possibility of fundamental particles and infinite virtual superpositions being one and the same.

The density of transverse modes of waves at frequency ω is $\omega^2 d\omega / \pi^2 c^3$ and the zero point energy for each of these modes is $\hbar\omega/2$. The electrostatic and magnetic energy densities in electromagnetic waves are equal, thus for electromagnetic zero point fields:

$$\frac{\overline{\varepsilon_0 E^2}}{2} + \frac{\overline{\varepsilon_0 c^2 B^2}}{2} = \frac{\hbar\omega}{2}\left[\frac{\omega^2 d\omega}{\pi^2 c^3}\right] \quad \text{and} \quad \overline{\varepsilon_0 E^2} = \overline{\varepsilon_0 c^2 B^2} = \frac{\hbar\omega^4}{2\pi^2 c^3}\frac{d\omega}{\omega}.$$

For a fundamental charge e using $\alpha = e^2 / 4\pi\varepsilon_0 \hbar c,$ and provided $\beta << 1,$ this gives an

$$\text{Average force squared of} \quad \overline{F^2} = e^2\overline{E^2} = \frac{2\alpha}{\pi}\frac{\hbar^2\omega^4}{c^2}\frac{d\omega}{\omega} \tag{2.1. 1}$$

Thinking semi classically, for an electron of rest mass m this can generate simple harmonic motion of amplitude r, where $F^2 = m^2\omega^4 r^2$ (if $\beta << 1$). Solving for r^2 (where r^2 is superimposed on the normal quantum mechanical electron orbit, $\lambdabar_c = \hbar / mc$ is the Compton

wavelength, and $k = \omega / c$):

$$r^2 = \frac{\hbar^2}{m^2c^2}\frac{2\alpha}{\pi}\frac{d\omega}{\omega} \quad = \quad \left[\lambdabar_c{}^2\right].\left[\frac{2\alpha}{\pi}\frac{dk}{k}\right]$$

Integrating r^2: $$r^2_{Total} = \lambda\!\!\bar{}_c{}^2 \frac{2\alpha}{\pi}\int_{k\min}^{k\max}\frac{dk}{k} \quad = \quad \lambda\!\!\bar{}_c{}^2 \frac{2\alpha}{\pi}\log(k_{\max}/k_{\min}).$$

The minimum and maximum values for k are chosen to fit atomic orbits, and a root mean square value for r can be found. Combining this with the small probability that the electron will be found in the nucleus, this small root mean square deviation shifts the average potential by approximately the Lamb shift. This can also be thought of as simple harmonic motion of amplitude $\approx \lambda\!\!\bar{}_c$, occurring with probability $(2\alpha/\pi)dk/k$. It can also be interpreted as the electron recoiling by $\approx \lambda\!\!\bar{}_c$ (when $\beta << 1$) in random directions due to virtual photon emission with a probability of $(2\alpha/\pi)dk/k$.

2.1.2 Dividing probabilities into the product of two component parts

This probability $(2\alpha/\pi)dk/k$ can be thought of as the product of two terms $A \& B$, where A includes the electromagnetic coupling constant α, B includes dk/k, and $AB = (2\alpha/\pi)dk/k$. This suggests that this same behaviour is possible if we have an appropriate superposition of virtual wavefunctions occurring with probability B, which emits virtual photons with probability A (by changing Eigenvalues $|\mathbf{p}_{nk}| = n\hbar k$ by $n = \pm 1$). For example, if a virtual superposition occurs with probability $B = (N \cdot s)dk/k$, and has a virtual photon emission probability for each member of these superpositions of $A = (N \cdot s)^{-1}(2\alpha/\pi)$, then the overall virtual photon emission probability remains as above at $AB = (2\alpha/\pi)dk/k$. This applies equally whether it is virtual gluon/photon/W&Z/graviton etc. emission. Provided A includes the appropriate coupling constant this same logic applies regardless of the type of boson emitted. As is usual to get integral or half integral total angular momentum $2s$ has to be integral and section 6.2 argues that the same applies for *N*. (This paragraph is simplified to illustrate the principle and will later be modified in section 3.3.)

In section 1.1.4 we said that these wavefunctions are built with squared vector potentials. If superpositions of them are to represent real particles they must be able to exist anywhere. This is possible only if they are generated by uniform fields. The only fields uniform in space-time are the zero point fields, and looking at the electromagnetic field first we can use section 2.1.1 above. Consider a vector $\mathbf{r}$ from some central origin O and a magnetic field vector $\mathbf{B}$ through origin O, then the vector potential at point $\mathbf{r}$ is $\mathbf{A} = (\mathbf{B} \times \mathbf{r})/2$ and the vector potential squared is $A^2 = (B^2 r^2 \sin^2\theta)/4$ where the angle between vectors $\mathbf{B} \& \mathbf{r}$ is θ.

As $\sin^2\theta$ averages 2/3 over a sphere: $\overline{A^2} = B^2 r^2 / 6$ (2.1. 2)

Here B^2 is the magnetic field squared at any point due to the cubic intensity of zero point EM also as in section 2.1.1. Putting Eq's. (2.1. 1) & (2.1. 2) together the vector potential squared is

$$\overline{e^2 A^2} = \frac{e^2 B^2 r^2}{6} = \frac{\alpha}{3\pi}\frac{\hbar^2\omega^4 r^2}{c^4}\frac{d\omega}{\omega} = \frac{\alpha}{3\pi}\hbar^2 k^4 r^2 \frac{dk}{k} \tag{2.1. 3}$$

As in section 2.1.2 we can divide this into two parts, noting the inclusion of spin *s* and integer *N* in the numerator and denominator:

$$\overline{e^2 A^2} = \left[\frac{\alpha}{3\pi sN}\hbar^2 k^4 r^2\right]\cdot\left[\frac{sN\cdot dk}{k}\right] \tag{2.1. 4}$$

But here a vector potential squared term $\left[\frac{\alpha}{3\pi sN}\hbar^2 k^4 r^2\right]$ *occurs with probability* $\left[\frac{sN\cdot dk}{k}\right]$.

Another way of looking at this is that a wavefunction ψ_k *that is generated by a vector potential squared term* $\left[\frac{\alpha}{3\pi sN}\hbar^2 k^4 r^2\right]$ *can occur with* $\left[\frac{sN\cdot dk}{k}\right]$ *probability.*

This is similar reasoning to that used in the semi classical Lamb shift explanation of section 2.1.1. In the first bracketed term of Eq. (2.1. 4), α is the electromagnetic coupling constant, but the same logic applies for the eight gluon and gravitational zero point vector fields where we will sum appropriate amplitudes of these and square this total as our effective coupling constant in Eq. (2.1. 4). But first we need to look at groups of spin zero preons that could build these wavefunctions. What mixtures of colours and electrical charges end up with the appropriate final colour and electrical charge for each of the fundamental particles or at least the ones we know of?

2.2 Spin Zero Virtual Preons from a Higgs type Scalar Field

2.2.1 Groups of eight preons that form superpositions

In this paper preons have zero spin and can have no weak charge. The only fields they can interact with (via Primary Interactions that build superpositions as in section 1.1.2) are colour, electromagnetic and gravity. In the simplest world there would be just one type of preon that comes in three colours, always positively charged say, with their three anti colours all negatively charged. We will assume that this is true unless it does not work. Looking at Table 2.2. 1 we see that a minimum of 6 preons is required to get the correct charge ratios of 3:2:1 between electrons, and up and down quarks. To get vector potential squared values that make all our equations work however, we need to couple to all 8 gluon fields requiring a total of 8 preons. Table 2.2. 1 has all the basic properties required to build infinite superpositions for the fundamental particles. We need to remember when looking at this table that from section 1.1.2 the effective secondary charge is much less than the primary charge and we have no idea yet of just what effective value the primary preon electric charge is.

Particles only are addressed in the groups of preons in Table 2.2. 1. To get anti particles it would seem that we can just change the signs of each preon in the groups of 8, excepting those that are already their own antiparticle. The first point to notice however is that both the electron and the W^- are predominantly anti preons, yet they are both defined as particles. Have we got something wrong? When we look at relativistic masses in section 3.2.1 we get the usual plus and minus solutions and Feynman showed us how to interpret the negative solutions as antiparticles. If this also applies in anti preons then because they are zero spin, and the weak force discriminates between particles and antiparticles by their helicity, this discrimination can apply only in secondary interactions. The preon anti preon content of the groups in Table 2.2. 1 does not necessarily tell us whether they produce particles or antiparticles. We will discuss this further in section 3.2.1, also as of now there is still no good understanding of the predominance of matter over antimatter in our universe. In Table 2.2. 1 only one example of colour is given for quarks and gluons. Different colours can be obtained by simply changing appropriate preon colours. Various combinations of 8 preons in this table are borrowed from a scalar field for time $\Delta T \approx \hbar / \Delta E$, this process continually repeating in time. Conservation of charge normally allows only opposite sign pairs of electric charges to appear out of the vacuum. Let us imagine that these virtual preons are building an electron for example whose electric charge exists continually unless it meets a positron and is annihilated.

Fundamental Particles	Preon colour	Preon electric charge.	Group colour	Group electric charge.
Spin 0 Higgs Spin ½ Neutrino family. Spin 1 Photons, Z_0 & Neutral gluons. Spin 2 Gravitons.	Any colour + its Anticolour Red Antired Green Antigreen Blue Antiblue	1 -1 1 -1 1 -1 1 -1	Colourless	0
Spin ½ Electron family. Spin 1 W^-.	Any colour + its Anticolour Antired Antired Antigreen Antigreen Antiblue Antiblue	1 -1 -1 -1 -1 -1 -1 -1	Colourless	-6
Spin ½ Blue up quark family.	Red Antired Green Antigreen Green Blue Blue Red	1 -1 1 -1 1 1 1 1	Blue	+4
Spin ½ Red down quark family.	Green Antigreen Red Antired Green Antigreen Antiblue Antigreen	1 -1 1 -1 1 -1 -1 -1	Red	-2
Spin 1 Red to green Gluons.	Red Antigreen Red Antired Green Antigreen Blue Antiblue	1 -1 1 -1 1 -1 1 -1	Red plus antigreen	0

Table 2.2. 1 Groups of 8 virtual preons forming the fundamental particles.

This charged electron is thus due to a continuous appearance out of and back into the vacuum of virtual charged preons in a steady state process existing for the life of the superposition, and not conflicting with conservation of charge. If the electron itself does not conflict then neither do the borrowed preons that build it.

2.2.2 Primary coupling constants behave differently

Q.E.D. tells us that the bare (electric) charge of an electron for example increases logarithmically inversely with radius from its centre. Polarizations of the vacuum (of virtual charged pairs) progressively shield the bare charge from a radius of approximately one Compton radius $\lambdabar_c$ inwards towards the centre. When an electron (for example) is created in some interaction the full bare charge is exposed for an infinitesimal time. Instantaneously after its creation, shielding due to polarization of the vacuum builds progressively outward from the centre of its creation at the velocity of light. For radii $\geq \lambdabar_c$ we measure the usual fundamental charge e. There are similar but more complicated processes that occur to the colour charge. Camouflage is the dominant one where the colour charge grows with radius as the emitted gluons themselves have color charge. At the instant of their birth the preons are bare and at this time $t = 0$ say, all the zero point vector fields can act on these bare colour and electric charges as there is simply no time for shielding and other effects to build. The primary coupling constants that we use must consequently be the same for all values of k in complete contrast to those for secondary interactions. We don't know what this primary electromagnetic coupling constant is so we will just call it α_{EMP}. Also we will find that to get any sense out of our equations the *primary colour coupling* has to be very close to 1. A coupling of 1 is a natural number and simply reflects certainty of coupling. Provided the *secondary colour coupling can be in line with the Standard Model* and there does not seem to be any other good reason to pick a number less than 1, we will make the (apparently arbitrary) *assumption that the bare primary colour coupling is exactly 1*. (In section 4.1.1 we will find that this seems to be consistent with the Standard Model.)

2.2.3 Primary interactions also behave differently

Let us define a frame in which the central origin of the wavefunctions ψ_k of our infinite superposition is at rest: The laboratory or rest frame we will refer to as the LF. The preons that build each ψ_k are born from a Higg's type scalar field with zero momentum in this frame. This has very relevant consequences as their *wavelength is infinite* in this rest frame at time $t = 0$, and after they build wavefunction ψ_k their wavelength is of the order $1/k$ for times

$0 < t < \hbar / 2E$. This implies that there could possibly be significant differences in the way amplitudes are handled between primary and secondary interactions.

Let us consider secondary interactions first with an electron and positron for example located approximately distance r apart. For photon wavelengths $<< r$ both the electron and the positron each emit virtual photons with probabilities proportional to α, but for wavelengths $>> r$ their amplitudes cancel. Returning to primary interactions, zero momentum preons must always have an infinite wavelength which is greater than the wavelengths (or $1/k$ values) of the zero point quanta they interact with, for all $k \neq 0$. This implies that we cannot simply add or subtract amplitudes algebraically; the charges are always further apart than the wavelength of the interacting quanta (except in the special case where $k = 0$). In fact if algebraic addition of amplitudes did apply in primary interactions, infinite superpositions for colourless and electrically neutral neutrinos would be impossible. So how can infinitely far apart preons of differing charge generate wavefunctions of all dimensions down to Planck scale? This can happen only if the amplitudes of all 8 preons are somehow linked over infinite space, all at the same time $t = 0$ contributing to generating the wavefunction ψ_k. This non-local behaviour is not new. Recent experiments have confirmed that what Einstein struggled to come to terms with is in fact true; he called it “spooky action at a distance”. While these experiments are so far limited in the distance over which they demonstrate entanglement, there is now wide acceptance that it can reach across the Universe. In the same manner wavefunctions covering all space can instantly collapse. We want to suggest here that this same non-locality applies in primary interactions: our 8 virtual preons all unite instantaneously at time $t = 0$ across infinite space in generating each ψ_k. Also the vector potential squared equations that they generate must always be the same for all the preon combinations in Table 2.2. 1. This can happen only if the amplitudes of all 8 are added regardless of charge sign for primary interactions. This applies to both colour and electric charge.

The opposite is true for the secondary interactions. At time $t = 0$ all 8 preons instantaneously collapse into some sort of virtual composite particle that for times $0 < t < \hbar / 2E$ obeys wavefunction ψ_k. The dimensions of ψ_k are of the same order as the wavelength of the interacting quanta, and the usual algebraic total electric charge and nett colour charge must now apply as in the group charges in Table 2.2. 1. All of this may seem contrary to current thinking which has gradually been built up over several centuries of secondary interaction experiments; however it may not be so out of place when viewed in the context of the counter intuitive results of entanglement experiments. The key point to bear in mind is that the predictions of this paper must agree or at least be able to fit the Standard Model, or secondary interaction experiments, as we may never be able to look into virtual primary interactions, but only observe their effects.

Amplitudes to interact are complex numbers which we can draw as a vector. This applies to both colour and electric coupling, where these two vectors can be at the same complex angle or at different angles. The simplest case is if they are in line and we will assume this is true for both colour and electromagnetic primary interactions which are both spin 1. This seems to work and when we later include gravity, a spin 2 interaction, we find that the spin 2 vector only works if it is at right angles to the two in line spin 1 vectors. *Let us start in a zero gravity world* by simply adding the 8 preon colour vectors of amplitude 1 and the eight primary electromagnetic vectors of amplitude $\sqrt{\alpha_{EMP}}$ together, as all this only works if they are all in line.

The total colour plus electromagnetic primary amplitude is $8+8\sqrt{\alpha_{EMP}}$ (2.2. 1)

This equation is always true regardless of signs as in section 2.2.3

The colour plus electromagnetic primary coupling constant is $\left(8+8\sqrt{\alpha_{EMP}}\right)^2$ (2.2. 2)

Inserting this into Eq. (2.1. 4) we get

$$Q^2A^2 = \left[\frac{\left[8+8\sqrt{\alpha_{EMP}}\right]^2}{3\pi sN}\hbar^2 k^4 r^2\right] \cdot \left[\frac{sN \cdot dk}{k}\right] \tag{2.2. 3}$$

Again we interpret this just as we did in section 2.1.2 and Eq. (2.1. 4) as a vector potential squared term

$$Q^2A^2 = \frac{\left[8+8\sqrt{\alpha_{EMP}}\right]^2}{3\pi sN}\hbar^2 k^4 r^2 \text{ occurring with probability } = \frac{sN \cdot dk}{k} \tag{2.2. 4}$$

Where Q is a symbol representing the entire 8 colour and 8 electric amplitudes combined, with s the spin and $N=2$ for infinitesimal mass infinite superpositions; but $N=1$ or $N=2$ for massive superpositions. (Table 4.3. 1, section 6 and its subsections cover this more fully.)

2.3 The Virtual Wavefunctions that form Infinite Superpositions

2.3.1 Infinite families of similar wavefunctions

Consider the family of wave functions where ignoring time:

$$\begin{aligned}\psi_{nk} &= U(nrk)Y(\theta\varphi) \\ U(nrk) &= C_{nk} r^l \exp(-n^2k^2r^2/18)\end{aligned} \tag{2.3. 1}$$

$U(nrk)$ is the radial part of ψ_{nk}, $Y(\theta\phi)$ the angular part, C_{nk} a normalizing constant, and we will find that l is the usual angular momentum quantum number. There is an infinite family of ψ_{nk}, one for each value k where $0 < k < \infty$ *in a zero gravity world.*

Now put $R(nrk) = rU(nrk) = C_{nk} r^{l+1} \exp(-n^2k^2r^2/18)$ (2.3. 2)

As we are dealing with zero spin preons we use Klein-Gordon equations [9]. The Klein-Gordon equation is based on the relativistic equation $\mathbf{p}^2 = E^2/c^2 - {m_0}^2c^2$ and in a squared vector potential the Time Independent Klein Gordon Equation is

$$\hat{P}^2\psi = -\hbar^2\nabla^2\psi + Q^2A^2\psi = \left[\frac{E^2}{c^2} - {m_0}^2c^2\right]\psi \tag{2.3. 3}$$

Using $\dfrac{\nabla^2\psi}{\psi} = \dfrac{1}{R}\dfrac{\partial^2 R}{\partial r^2} - \dfrac{l(l+1)}{r^2}$ we get the Time Independent

Radial Klein Gordon Equation $\dfrac{\hbar^2}{R}\dfrac{\partial^2 R}{\partial r^2} = \dfrac{l(l+1)\hbar^2}{r^2} + Q^2A^2 - \left[\dfrac{E^2}{c^2} - m_0c^2\right]$ (2.3. 4)

For each ψ_{nk} the energy is E_{nk} a function of $n \,\&\, k$, and we will label the rest mass as m_{0snk} a function of spin s, $n \,\&\, k$, but also a function of the particle rest mass m_0 and this becomes

$$\frac{\hbar^2}{R}\frac{\partial^2 R}{\partial r^2} = \frac{l(l+1)\hbar^2}{r^2} + Q^2A^2 - \left[\frac{{E_{nk}}^2}{c^2} - {m_{0snk}}^2c^2\right] \tag{2.3. 5}$$

Differentiating $R(nrk) = rU(nrk) = C_{nk} r^{l+1} \exp(\frac{-n^2k^2r^2}{18})$ twice with respect to r, multiplying by $\hbar^2$ and dividing by R

$$\frac{\hbar^2}{R}\frac{\partial^2 R}{\partial r^2} = \frac{l(l+1)\hbar^2}{r^2} + \frac{n^4\hbar^2k^4r^2}{81} - \frac{(2l+3)n^2\hbar^2k^2}{9} \qquad (2.3.\ 6)$$

Comparing Eq's. (2.3. 5) & (2.3. 6) we see that l is the usual angular momentum quantum number and the vector potential squared required to generate these wavefunctions is

$$Q^2A^2 = \frac{n^4\hbar^2k^4r^2}{81} = \left[\frac{n}{3}\right]^4 \hbar^2k^4r^2 \qquad (2.3.\ 7)$$

The momentum squared is $\mathbf{p}_{nk}{}^2 = \frac{E_{nk}{}^2}{c^2} - m_{0snk}{}^2c^2 = \frac{(2l+3)n^2\hbar^2k^2}{9}$ (2.3. 8)

For $l = 3$ wavefunctions this becomes $\mathbf{p}_{nk}{}^2 = n^2\hbar^2k^2$ & $|\mathbf{p}_{nk}| = n\hbar k$ (2.3. 9)

2.3.2 Eigenvalues of these wavefunctions and parallel momentum vectors

From Eq.'s (2.3. 8) & (2.3. 9) as $k \to \infty$, the energy squared $E_{nk}{}^2 \to \mathbf{p}_{nk}{}^2c^2 = n^2\hbar^2\omega^2$ or

If $l = 3$ when $k \to \infty$ energy $E_{nk} \to n\hbar\omega$ (taking the positive case). (2.3. 10)

In virtual exchange processes $m_{0X}{}^2c^4 = E_X{}^2 - \mathbf{p}_X{}^2c^2$ (where subscript X refers to the virtual exchanged quantities) is not always true. Also there is no mass exchange here as the preons are born with a virtual rest mass at $t = 0$, followed instantaneously by this interaction. As $k \to \infty$ however, the virtual quanta exchanged start to behave as real; each quanta transferring energy $\hbar\omega$, and total energy $n\hbar\omega$, if n is integral. When $k << \infty$ however, the energy transferred $E_X{}^2 < \mathbf{p}_{nk}{}^2c^2$ and $< n^2\hbar^2\omega^2$, but as virtual quanta can transfer $|\mathbf{p}| = \hbar k$ and at the same time have $E_X < \hbar\omega$, this allows n to remain integral for all k. This can be thought of as n parallel momentum vector $|\mathbf{p}| = \hbar k$ quanta, transferring total momentum $|\mathbf{p}_{nk}| = n\hbar k$ and energy $E_X \le n\hbar\omega$ to wavefunction ψ_{nk}. This implies $\psi_{nk} = C_{nk}r^3\exp(-n^2k^2r^2/18)$ wavefunctions have (where $0 < k < \infty$ *until we introduce gravity* when $k_{Cutoff} < \infty$)

Eigenvalues $\mathbf{p}_{nk}{}^2 = n^2\hbar^2k^2$ if n is integral and $l = 3$. (2.3. 11)

Also there are no scalar potentials involved, only squared vector potentials, so this is a magnetic or vector type interaction. Particles in classical magnetic fields have a constant magnitude of linear momentum which is consistent with the squared momentum Eigenvalues of Eq. (2.3. 11).*This also implies that each* ψ_{nk} *is formed from quanta of wave number* k *only and that secondary interactions with* ψ_{nk} *emit or absorb* $|\hbar k|$ *virtual quanta if* n *changes by* ± 1. The wavefunction ψ_{nk} is virtual and in this sense both the energy E_{nk} and rest mass m_{0snk} in Eq. (2.3. 8) are also virtual quantities borrowed from some sort of scalar Higg's field. We use these virtual quantities to calculate the amplitude that ψ_{nk} is in an m state of angular momentum in section 3.1, and in section 3.2 to calculate total angular momentum and real rest mass. As in section 2.3.2 above, we can think of $|\mathbf{p}_{nk}| = n\hbar k$ as n parallel momentum vectors $|\mathbf{p}| = \hbar k$. As spin 3 (or $l = 3$) needs at least 3 spin 1 quanta to build it n must be at least 3. When $n = 3$ we can think of this as 3 of the 8 preons each absorbing quanta $|\hbar k|$ at time $t = 0$. We will find that a spin ½ state has a dominant $n = 6$ mode where 6 of the 8 preons each absorb quanta $|\hbar k|$. It needs at least two smaller side modes $n = 5$ & $n = 7$ with either 5 or 7 of the 8 preons each absorbing quanta $|\hbar k|$ respectively at $t = 0$ as in Figure 3.1. 4 where a positron superposition is illustrated.

From Eq. (2.3. 7) $Q^2A^2 = \left[\dfrac{n}{3}\right]^4 \hbar^2 k^4 r^2 = 16\hbar^2 k^4 r^2$ for this dominant $n = 6$ mode.

Thus from Eq. (2.2. 4) $Q^2A^2 = \dfrac{\left[8 + 8\sqrt{\alpha_{EMP}}\right]^2}{3\pi s} \hbar^2 k^4 r^2 = 16\hbar^2 k^4 r^2$ for an $n = 6$ mode.

Now s = ½ for spin ½ and $\dfrac{2\left[8 + 8\sqrt{\alpha_{EMP}}\right]^2}{3\pi} = 16$ if we have only an $n = 6$ mode.

Thus $8 + 8\sqrt{\alpha_{EMP}} = \sqrt{24\pi}$ and $\alpha_{EMP}^{-1} \approx 137.1$, but this is true for an $n = 6$ mode only, and we have a superposition where the amplitudes of the smaller side modes $n = 5$ & $n = 7$ determine the ratio between the primary to secondary (colour and electromagnetic) coupling amplitudes or the value of $\alpha_3^{-1} @ L_P$ (Section 3.3). The Q^2A^2 required to produce this superposition with amplitudes c_n is, using Eq. (2.3. 7)

$$Q^2A^2 = \sum_{n=5,6,7} c_n * c_n \frac{n^4 \hbar^2 k^4 r^2}{81} \tag{2.3. 12}$$

Repeating the same procedure as above for these superpositions and using Eq. (2.3. 12) we find the strength of α_{EMP} required increases considerably (see sections 4.1.1 & 4.3). As the secondary electromagnetic coupling $\alpha_{EMS}{}^{-1} @ L_P$ must be constant for all spin ½ leptons and quarks the amplitudes of the smaller side modes $n=5$ & $n=7$ that determine this must also be constant for all the fermions, implying that Eq. (2.3. 12) must be the same for all fermions. The same arguments apply to the other groups of fundamental particles but we will return to this in sections 3.3 where we see that the same also applies with graviton emission.

3 Properties of Infinite Superpositions

3.1 What is the Amplitude that ψ_{nk} is in an *m* state?

3.1.1 Four vector transformations

The rules of quantum mechanics tell us that if we carry out any measurement on a real spherically symmetric $l=3$ wavefunction it will immediately fall into one of the seven possible states $l=3, m=0,\pm1,\pm2,\pm3$ [10]. But ψ_{nk} is a virtual $l=3$ wave function so we cannot measure its angular momentum. During its brief existence it must always remain in some virtual superposition of the above seven possible states and we can describe only the amplitudes of these. So is there any way to calculate these amplitudes as they must relate to the amplitudes of the angular momentum states of the spin 1 quanta it absorbs from the zero point vector fields? First consider the 4 vector wavefunction of a spin 1 particle and start with a time polarized state which has equal probability of polarization directions. It is thus spherically symmetric, which we will label as ss. Using 4 vector (*t, x, y, z*) notation:

In frame A, a time polarized or ss spin 1 state is (1,0,0,0).

Let frame B move along the z axis at velocity $\beta = v/c$ in the z direction.

In frame B the polarization state transforms to $(\gamma, 0, 0, \gamma\beta)$.

But this is γ^2 time polarized (ss states) minus $\gamma^2\beta^2 \times z$ polarized ($m=0$ states).

In frame B there are $\gamma^2 \times ss$ states $-\gamma^2\beta^2 \times m=0$ states.

Now $\gamma^2 - \gamma^2\beta^2 = \gamma^2(1-\beta^2) = 1$ is an invariant probability in all frames and in removing $\gamma^2\beta^2 \times (m=0)$ states from γ^2 ss states, the new ratio of spherical symmetry is $(\gamma^2 - \gamma^2\beta^2)/\gamma^2 = 1-\beta^2$. Thus a spherically symmetric state is transformed from probability

1 in frame A, to $1-\beta^2$ in frame B. Also removing $m=0$ states from spherically symmetric states leaves a surplus of $m=\pm 1$ states, as spherically symmetric states are equal superpositions of $m=-1,\ m=0, \& \, m=+1$ states.

Thus in Frame B the probabilities are $(1-\beta^2)\times ss$ states $+\beta^2 \times m=\pm 1$ states. (3.1. 1)

We can describe this as a virtual superposition of $(\frac{1}{\gamma}\times ss, +\beta\times m=\pm 1)$ states. (3.1. 2)

As $\beta^2 \to 1$ we have transverse polarized states, the same as real photons. Now transverse polarized spin 1 states can be either left $(m=-1)$, or right $(m=+1)$ circular polarization, or equal superpositions of $(1/\sqrt{2})L+(1/\sqrt{2})R$ as in $x \,\&\, y$ polarization. If we think of individual spin zero preons absorbing these spin 1 quanta at $t=0$ they must also have this same β^2 probability of transversely polarized spin 1 states. If they then merge into some composite $l=3$ particle (as in Figure 3.1. 4) for time $0<t<\hbar/2E,$ the probability of it being in some particular state $(l=3, m=0),\ (l=3, m=\pm 1),\ (l=3, m=\pm 2)$ or $(l=3, m=\pm 3)$, must be the same β^2. If we look at Eq.'s (1.1. 1) we can see what is behind them. We have written the amplitudes in these three equations in terms of β_{nk} & γ_{nk} as this is the most convenient way to express them. Velocities however are not the quantum world language, momenta are. Velocity operators are momentum operators over relativistic masses. Our Eigenvalues are $\mathbf{p}_{nk}{}^2 = n^2\hbar^2k^2$ for each $n \,\&\, k$, and this allows our velocity operators to give constant $\beta_{nk}{}^2$. Even though the mass in these operators is virtual, we can still use it to calculate $|\beta_{nk}|$. For each k and integral n there will be a constant $|\beta_{nk}|$ and $\gamma_{nk} = (1-\beta_{nk}{}^2)^{-1/2}$. As we will see β_{nk} can be thought of as the magnitude of the velocity of an imaginary centre of momentum frame in which these interactions take place. We will also draw our Feynman diagrams of these interactions in terms of $\beta_{nk} \,\&\, \gamma_{nk}$ for convenience, even though this is unconventional. To proceed from here however, we need to define two frames as follows:

1) The Laboratory Frame (LF) or Fixed Frame as in section 2.2.3

The infinite superposition has rest mass m_0 and zero nett momentum in this frame. Each ψ_{nk} is centered here with magnitude of momentum $|\mathbf{p}_{nk}| = n\hbar k$. Even though we have no idea of the direction of this momentum vector we will define it as the z direction. The eight preons are born in this frame with zero momentum and can thus be considered here as being at rest or with zero velocity and infinite wavelength at their birth. The Feynman diagram of the interaction in this frame that builds ψ_{nk} is illustrated in Figure 3.1. 3.

2) The Center of Momentum Frame (CMF)

This (imaginary) frame is the center of momentum of the interaction that builds ψ_{nk}. The CMF moves at velocity β_{nk} relative to the laboratory frame in the z direction or parallel to the unknown momentum vector direction $\mathbf{p}_{nk}$. In this CMF the momenta and velocities of the preons at birth and after the interaction are equal and opposite. This is illustrated in Figure 3.1. 2 again in terms of $m_0, \beta_{nk}, \& \gamma_{nk}$. In the LF the velocity of the preons at birth is zero, in the CMF this is $-\beta_{nk}$ and after the interaction $+\beta_{nk}$, where both $-\beta_{nk}$ and $+\beta_{nk}$ are in the unknown z direction. In the LF the particle velocity $\beta_{nk}(particle) = \beta_{nkp}$ is the simple relativistic addition of the two equal velocities β_{nk} as in Figure 3.1. 1.

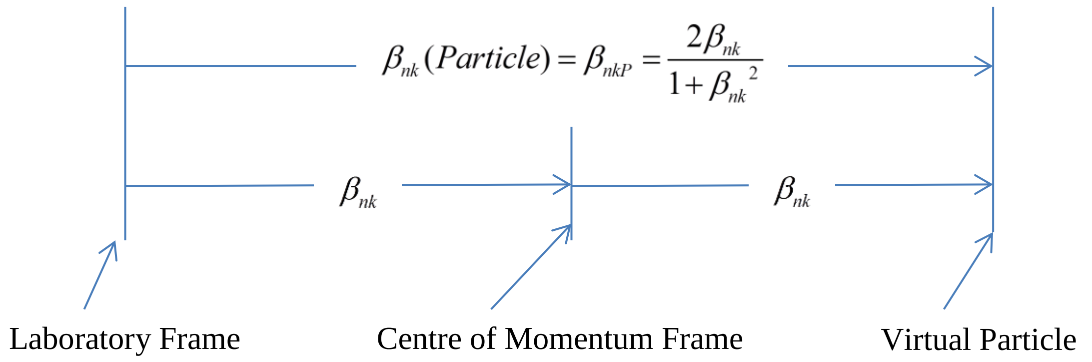


Figure 3.1. 1

3.1.2 Feynman diagrams of primary interactions

Let us start with

$$\beta_{nk}(Particle) = \beta_{nkP} = \frac{2\beta_{nk}}{1+\beta_{nk}^2} \text{ and } \gamma_{nkP} = (1-\beta_{nkp}^2)^{-1/2} = \gamma_{nk}^2(1+\beta_{nk}^2) \tag{3.1. 3}$$

If the particle rest mass is m_0 let each preon have a virtual rest mass $m_0/(8\gamma_{nk}\sqrt{2s})$.

The eight preons are effectively a virtual particle of rest mass $m_{0snk} = \dfrac{m_0}{\gamma_{nk}\sqrt{2s}}$ (3.1. 4)

The particle momentum in the LF is zero at birth. After the interaction using these equations

$$|\mathbf{p}_{nk}| = n\hbar k = m_{0snk}\beta_{nkP}\gamma_{nkP}c = \left[\frac{m_0}{\gamma_{nk}\sqrt{2s}}\right]\left[\frac{2\beta_{nk}}{1+\beta_{nk}^2}\right]\left[\gamma_{nk}^2(1+\beta_{nk}^2)\right]c$$

The particle momentum after the interaction in the LF $|\mathbf{p}_{nk}| = n\hbar k = \dfrac{2m_0\beta_{nk}\gamma_{nk}c}{\sqrt{2s}}$ (3.1. 5)

Using Eq. (3.1. 4), in the LF the particle energy at birth is

$$m_{0snk}c^2 = \frac{m_0c^2}{\gamma_{nk}\sqrt{2s}} \tag{3.1. 6}$$

In the LF the particle energy after the interaction is using Eq's. (3.1. 3)

$$m_{0snk}\gamma_{pnk}c^2 = \frac{m_0}{\gamma_{nk}\sqrt{2s}}\gamma_{nk}{}^2(1+\beta_{nk}{}^2)c^2 = \frac{m_0\gamma_{nk}}{\sqrt{2s}}(1+\beta_{nk}{}^2)c^2 \tag{3.1. 7}$$

In the CMF the momentum at birth is using Eq. (3.1. 4)

$$-m_{0snk}\gamma_{nk}\beta_{nk} = \frac{-m_0\beta_{nk}}{\sqrt{2s}} \tag{3.1. 8}$$

In the CMF the momentum after the interaction is equal but in the opposite direction

$$= \frac{+m_0\beta_{nk}}{\sqrt{2s}} \tag{3.1. 9}$$

In the CMF the energy at birth, and after the interaction is

$$m_{0snk}\gamma_{nk}c^2 = \frac{m_0c^2}{\sqrt{2s}} \tag{3.1. 10}$$

These values are all summarized in Figure 3.1. 2 and Figure 3.1. 3 but with $c = 1$.

From Eq. (3.1. 5) $|\mathbf{p}_{nk}| = n\hbar k = \dfrac{2m_0\beta_{nk}\gamma_{nk}c}{\sqrt{2s}}$ and $\beta_{nk}\gamma_{nk} = \dfrac{n\hbar k\sqrt{2s}}{2m_0c} = \dfrac{\lambda\!\!\!^{-}_c nk\sqrt{2s}}{2}$

(where $\lambda\!\!\!^{-}_c$ is the Compton wavelength) so let us now put:

$$K_{nk} = \beta_{nk}\gamma_{nk} = \frac{n\hbar k\sqrt{2s}}{2m_0c} = \frac{\lambda\!\!\!^{-}_c nk\sqrt{2s}}{2} \tag{3.1. 11}$$

In terms of K_{nk}: $\beta_{nk}{}^2 = \dfrac{K_{nk}{}^2}{1+K_{nk}{}^2}$ and $\gamma_{nk}{}^2 = 1+K_{nk}{}^2$ (3.1. 12)

Each infinite superposition has a fixed $\lambda\!\!\!^{-}_c$, also each wavefunction ψ_{nk} of this infinite superposition has fixed n & s, thus $K_{nk} \propto k$. For example we can put

$$\frac{dK_{nk}}{K_{nk}} = \frac{dk}{k} \tag{3.1. 13}$$

These simple expressions and what follows are not possible if $m_{0snk} \neq m_0 / \gamma_{nk}\sqrt{2s}$, and *when we include gravity* we find $m_{0snk} = m_0 / (\gamma_{nk}\sqrt{2s})$ is essential (section 4.2).

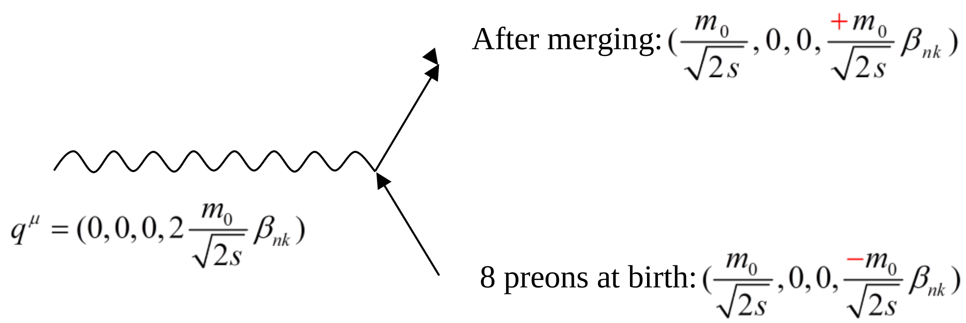


Figure 3.1. 2 Feynman diagram in an imaginary centre of momentum frame.

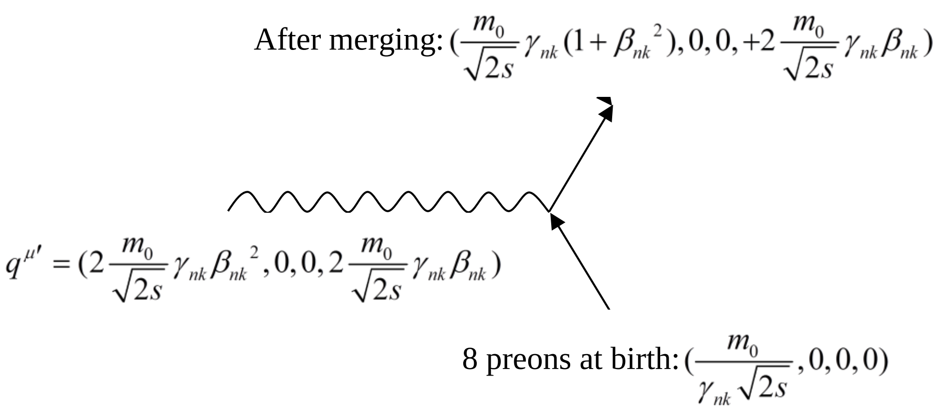


Figure 3.1. 3 Feynman diagram in the laboratory frame.

The interaction in the Feynman diagrams above is with spin 1 quanta. The Feynman transition amplitude of this interaction tells us that the polarization states of these exchanged quanta is determined by the sum of the components of the initial, plus final 4 momentum $(p_i + p_f)^{\mu}$. Ignoring all other common factors this tells us that the space polarized component is the sum of the momentum terms $(\mathbf{p}_i + \mathbf{p}_f)$ and the time polarized component is the sum of the energy terms $(p_i + p_f)^0$. We have defined our momentum as in an unknown z direction:

The ratio of z polarization to time polarization amplitudes is $\dfrac{(p_i + p_f)^z}{(p_i + p_f)^0}$ (3.1. 14)

In the CMF $(p_i + p_f)^z = 0$, thus an interaction in the CMF exchanges only time polarized, or spherically symmetric $l = 1$ states. In the LF the ratio of z (or $m = 0$) polarization, to time polarization in the LF is $\beta_{nk}{}^2$,

$$\text{where} \quad \beta_{nk} = \frac{(p_i + p_f)^z}{(p_i + p_f)^0} = \frac{2m_0\gamma_{nk}\beta_{nk}}{2m_0\gamma_{nk}} \qquad (3.1.\ 15)$$

From section 3.1.1 these are probabilities of $\gamma_{nk}{}^2\ ss\ -\gamma_{nk}{}^2\beta_{nk}{}^2 \times (m = 0)$ states, or $(1-\beta_{nk}{}^2)\ ss\ +\beta_{nk}{}^2 \times (l = 1, m = \pm 1)$ states.

In the LF this is a virtual superposition of $(\frac{1}{\gamma_{nk}} \times ss, +\beta_{nk} \times m = \pm 1)$ states. (3.1. 16)

From section 3.1.1 as these quanta from the scalar and vector zero point fields build each ψ_{nk} this implies that:

In the LF ψ_{nk} has virtual superposition amplitudes $(\frac{1}{\gamma_{nk}} \times ss, +\beta_{nk} \times m$ states) (3.1. 17)

From section 3.1.1 appropriate $l = 1, m = \pm 1$ superpositions can build any $l = 3,\ m$ state. Figure 3.1. 4 is an example of such a ψ_{nk} for $n = 5, 6, \&7\ (l = 3, m = +2)$ states.

3.1.3 Different ways to express superpositions

We have expressed all superpositions here in terms of spherically symmetric and m states for convenience and simplicity. We could have expressed them in the form:

$$\frac{1}{\gamma_{nk}\sqrt{7}}\left[(m=-3),(m=-2),(m=-1),(m=0),(m=+1),m=+2),(m=+3)\right] + \beta_{nk}(m=+2)$$

This is equivalent to (ignoring complex number amplitude factors)

$$\psi_{nk} = \frac{1}{\gamma_{nk}} \times ss, +\beta_{nk} \times (m=+2) \text{ where we have put m} = +2 \text{ for example.}$$

Because all these wavefunctions are virtual they cannot be measured in the normal way that collapses them into any of these Eigenstates, it is more convenient to use the method adopted here which is similar to QED virtual photons superpositions.

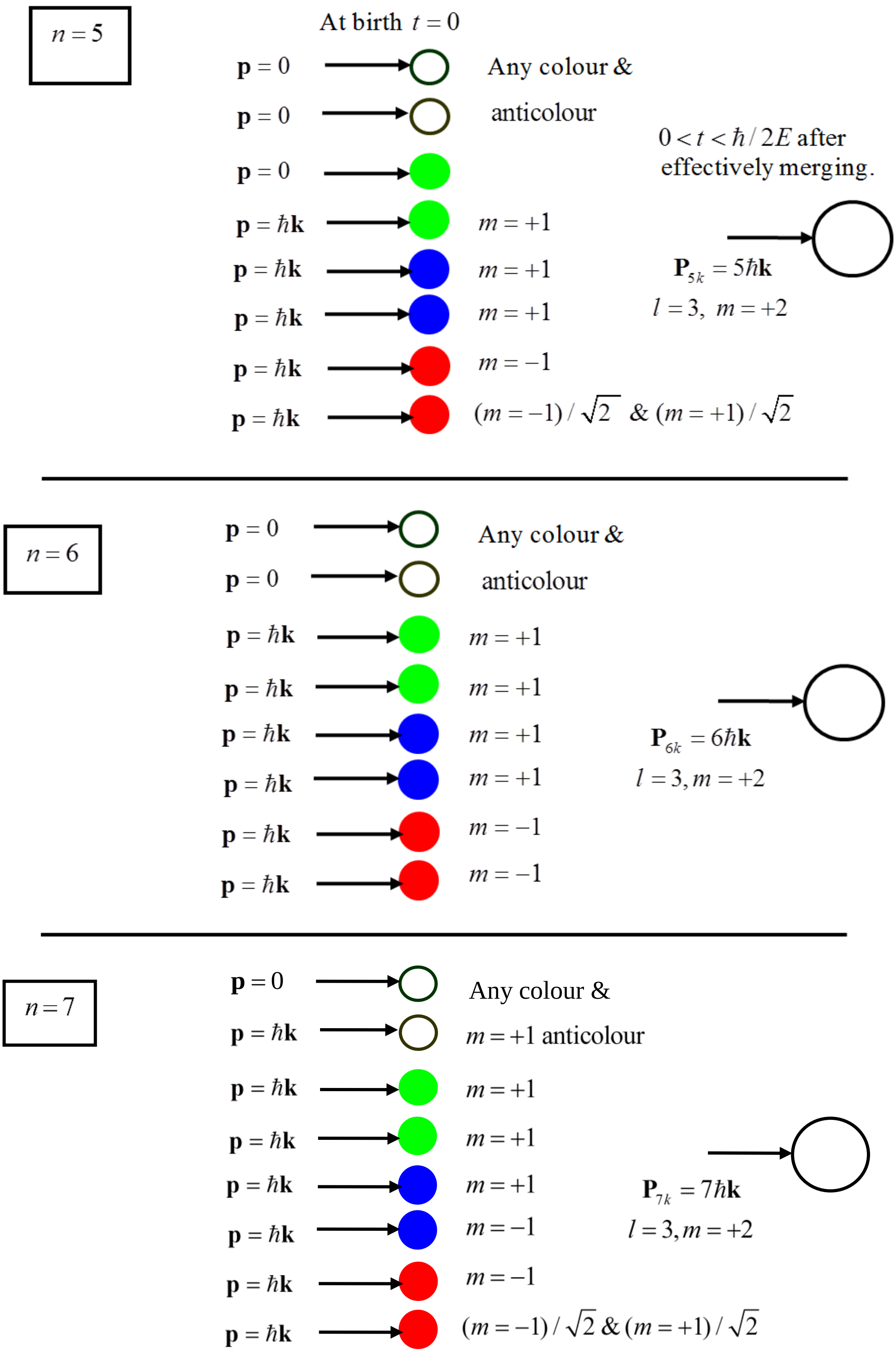


Figure 3.1. 4 Eight preons forming $m=+2$ states as part of a positron superposition.

There is no significance in which preons absorb quanta in the above.

3.2 Mass and Total Angular Momentum of Infinite Superpositions

3.2.1 Total mass of massive infinite superpositions

We will consider the total mass first of the infinite assembly and consider one wavefunction ψ_{nk} at a time, temporarily assuming that the amplitude c_n of each ψ_{nk} has magnitude $|c_n| = 1$. Each time ψ_{nk} is born it borrows virtual mass and virtual energy from scalar and vector zero point fields. We are however going to focus solely on the Eigenvalues $\mathbf{p}_{nk}{}^2 = n^2\hbar^2k^2$, and treat only momentum as real. But do these momenta themselves also leave momentum debts in the vacuum?

From section 2.3.2 as $k \to \infty$, $E_{nk}{}^2 \to \mathbf{p}_{nk}{}^2c^2 = n^2\hbar^2\omega^2$ or $E_{nk} \to n\hbar\omega$ and n quanta of energy $\hbar\omega$ and momentum $|\hbar k|$ are absorbed. We know that in some unknown direction $\mathbf{p}_{nk} = n\hbar\mathbf{k}$, which implies these n absorbed quanta must leave a cancelling debt in the opposite direction of $\mathbf{p}_{nk}(debt) = -n\hbar\mathbf{k}$ in the vacuum. But this is true only as $k \to \infty$ & $\beta_{nk}{}^2 \to 1$ and the virtual quanta energy transferred $E_X \to \hbar\omega$. So what happens when $\beta_{nk}{}^2 << 1$? Our wavefunctions ψ_{nk} are generated from a vector potential squared term A^2 derived in section 2.1.2 which in turn came from the B^2 term of section 2.1.1. As discussed in section 2.3.2 the Eigenvalues $\mathbf{p}_{nk}{}^2 = n^2\hbar^2k^2$ confirm the constant momentum squared feature of magnetic interactions. Also in section 2.1.1 the scalar virtual photon emission probability is directly related to the force squared term $F^2 = \varepsilon^2E^2$. Magnetic coupling probabilities are related to a magnetic force squared term $F^2 = \beta^2\varepsilon^2B^2/c^2 = \beta^2\varepsilon^2E^2$, where from section 3.1.2 and Eq. (3.1. 14) the ratio of this scalar to magnetic coupling is $\beta_{nk}{}^2$. Thus when $k < \infty$ and the exchanged energy $E_X \neq \hbar\omega$, $\beta_{nk}{}^2n$ quanta $|\hbar k|$ are absorbed from the vacuum and:

$$\text{We can expect a momentum debt of } \mathbf{p}_{nk}(debt) = -\beta_{nk}{}^2n\hbar\mathbf{k} \tag{3.2. 1}$$

We can now take the next step, as both the vectors $\mathbf{p}_{nk}$ and $\mathbf{p}_{nk}(debt)$ are parallel in the same unknown direction, to get a nett momentum:

$$\mathbf{p}_{nk}(nett) = \mathbf{p}_{nk} + \mathbf{p}_{nk}(debt) = (1-\beta_{nk}{}^2)n\hbar\mathbf{k} = \frac{n\hbar\mathbf{k}}{\gamma_{nk}{}^2} = \frac{\mathbf{p}_{nk}}{\gamma_{nk}{}^2} \text{ at wavenumber } k. \tag{3.2. 2}$$

Does this make sense using the relativistic energy expression $E_n{}^2 = \sum_{k=0}^{k=\infty}\mathbf{p}_{nk}(nett)^2c^2$ as we have simply ignored all other borrowed virtual energy and virtual mass?

We will initially look at only $N=1$ in Eq.(2.2. 4) massive infinite superpositions. Thus using probability $sN\cdot dk/k = s\cdot dk/k$, also Eq's.(3.1. 11), (3.1. 12),(3.1. 13),&(3.2. 2).

$$E_n{}^2 = c^2\int_{k=0}^{k=\infty}\mathbf{p}_{nk}(nett)^2\,\frac{s\cdot dk}{k} = c^2\int_0^\infty \frac{n^2\hbar^2k^2}{\gamma_{nk}{}^4}\frac{s\cdot dk}{k} = 4m_0{}^2c^4\int_0^\infty \frac{K_{nk}{}^2}{(1+K_{nk}{}^2)^2}\frac{dK_{nk}}{2K_{nk}}$$

$$E_n{}^2 = m_0{}^2c^4\left[\frac{-1}{1+K_{nk}{}^2}\right]_0^\infty = m_0{}^2c^4 \text{ or } E_n = \pm m_0c^2 \tag{3.2. 3}$$

This energy is due to summing momenta squared but it is real, with a real mass $\pm m_0$ for infinite superpositions of wavefunctions ψ_{nk}. These superpositions can form all the non infinitesimal mass fundamental particles. The equations do not work if the mass m_0 is zero. (We will look at infinitesimal masses in section 6.2.) The negative mass solutions in Eq. (3.2. 3) must be handled in the usual Feynman manner, and treated as antiparticles with positive energy going backwards in time. If they are spin ½ this also determines how they interact with the weak force.

3.2.2 Angular momentum of massive infinite superpositions

We will use the same procedure for the total angular momentum of $N=1$ in Eq.(2.2. 4) type infinite superposition with non infinitesimal mass.

Wavefunctions $\psi_{nk} = C_{nk}r^3\exp(-n^2k^2r^2/18)Y(\theta,\varphi)$ have angular momentum squared Eigenvalues $\mathbf{L}^2 = 12\hbar^2$ and m states of this have angular momentum Eigenvalues $\mathbf{L}_z = m\hbar$, thus we will treat both angular momentum and angular momentum debts as real just as we did for linear momentum. Even though m state wavefunctions are part of superpositions they still have probabilities just as the linear momenta squared above and it seemed to work. Using exactly the same arguments as in section 3.2.1 , if ψ_{nk} is in a state of angular momentum $\mathbf{L}_{zk} = m\hbar$, then it must leave an angular momentum debt in the vacuum of $\mathbf{L}_{zk}(debt) = -\beta_{nk}{}^2 m\hbar$ (or as in section) $\mathbf{L}_{zk}(nett) = \mathbf{L}_{zk} - \mathbf{L}_{zk}(debt)$.

$$\mathbf{L}_{zk}(nett) = (1-\beta_{nk}{}^2)m\hbar = (1-\beta_{nk}{}^2)\mathbf{L}_{zk} = \frac{\mathbf{L}_{zk}}{\gamma_{nk}{}^2} \quad (\text{if } \mathbf{L}_{zk} = m\hbar) \tag{3.2. 4}$$

But from Eq. (3.1. 16) the probability that $\mathbf{L}_{zk}$ is in an m state is also $\beta_{nk}{}^2$ so that

$$\text{Including } \beta_{nk}{}^2 \text{ probability } \mathbf{L}_{zk}(nett) = m\hbar\frac{\beta_{nk}{}^2}{\gamma_{nk}{}^2} \text{ at wavenumber } k. \tag{3.2. 5}$$

For an $N=1$ type infinite superposition $\mathbf{L}_z(Total) = \int_{k=0}^{k=\infty} \mathbf{L}_{zk}(nett)\dfrac{s\cdot dk}{k}$. $= sm\hbar\int_0^\infty \dfrac{\beta_{nk}{}^2}{\gamma_{nk}{}^2}\dfrac{dk}{2k}$

Using Eq's. (3.1. 11) to (3.1. 13) $\mathbf{L}_z(Total) = sm\hbar\int_0^\infty \dfrac{K_{nk}{}^2}{(1+K_{nk}{}^2)^2}\dfrac{dK_{nk}}{K_{nk}} = \dfrac{sm\hbar}{2}\left[\dfrac{-1}{1+K_{nk}{}^2}\right]_0^\infty$

$$\mathbf{L}_z(Total) = m'\hbar = \frac{sm\hbar}{2} \qquad \text{or} \qquad m' = \frac{s}{2}m \tag{3.2. 6}$$

Where m' *is the angular momentum state of the infinite superposition and* m *the state of* ψ_{nk} *.* Thus for spin ½ particles with $s=$ ½ in Eq.(3.2. 6) $m' = m/4$, but m' can be only ± ½, implying that the m state of ψ_{nk} that generates spin ½ must be $m = \pm 2$. An $N=1$ massive spin 1 particle has $s=1$ with $m' = m/2$. ($N=2$ is covered in section 6.2.) This is summarized in the following three member infinite superpositions.

$$(N=1)\ \text{Spin } ½ \qquad \psi_{\infty,1/2,\pm 1/2} = \sum_{n=5,6,7} c_n \int_{k=0}^{k=\infty}\left[\frac{1}{\gamma_{nk}}(\psi_{nk,ss}) + \beta_{nk}(\psi_{nk,\pm 2})\right]\sqrt{\frac{1}{2k}}dk \tag{3.2. 7}$$

$$(N=1)\ \ \text{Spin } 1 \qquad \psi_{\infty,1,m} = \sum_{n=4,5,6} c_n \int_{k=0}^{k=\infty}\left[\frac{1}{\gamma_{nk}}(\psi_{nk,ss}) + \beta_{nk}(\psi_{nk,2m})\right]\sqrt{\frac{1}{k}}dk \tag{3.2. 8}$$

The spin vectors of each ψ_{nk} with $|\mathbf{L}| = 2\sqrt{3}\hbar$, and their spin vector debts in the zero point vector fields, have to be aligned such that the sum in each case is the correct value: $|\mathbf{L}| = \sqrt{3}\hbar/2$, $|\mathbf{L}| = \sqrt{2}\hbar$ or $|\mathbf{L}| = \sqrt{6}\hbar$ for spins ½ , 1 & 2 respectively. Gravity (the ε term) is included in Eq. (1.1. 1) in our summary also spin 1 in Eq. (3.2. 8) is for $N=1$.

Spherically symmetric massive $N=1$ spin 1 states are superpositions of the three states $\dfrac{1}{\sqrt{3}}(m'=-1), \dfrac{1}{\sqrt{3}}(m'=0), \dfrac{1}{\sqrt{3}}(m'=+1)$, and using Eq. (3.2. 8) can be formed as follows

$$\text{Massive spin 1} \quad \left[\begin{aligned} &\frac{1}{\sqrt{3}}\psi_{\infty,1,m'=-1} = \frac{1}{\sqrt{3}}\sum_{n=4,5,6} c_n \int_{k=0}^{k=\infty}\left[\frac{1}{\gamma_{nk}}(\psi_{nk,ss}) + \beta_{nk}(\psi_{nk,m=-2})\right]\sqrt{\frac{1}{k}}dk \\ +&\frac{1}{\sqrt{3}}\psi_{\infty,1,m'=0} = \frac{1}{\sqrt{3}}\sum_{n=4,5,6} c_n \int_{k=0}^{k=\infty}\left[\frac{1}{\gamma_{nk}}(\psi_{nk,ss}) + \beta_{nk}(\psi_{nk,m=0})\right]\sqrt{\frac{1}{k}}dk \\ +&\frac{1}{\sqrt{3}}\psi_{\infty,1,m'=+1} = \frac{1}{\sqrt{3}}\sum_{n=4,5,6} c_n \int_{k=0}^{k=\infty}\left[\frac{1}{\gamma_{nk}}(\psi_{nk,ss}) + \beta_{nk}(\psi_{nk,m=+2})\right]\sqrt{\frac{1}{k}}dk \end{aligned}\right] \tag{3.2. 9}$$

3.2.3 Mass and angular momentum of complete infinite superpositions of ψ_k

In sections 3.2.1 & 3.2.2 to keep things simple we looked only at one wavefunction ψ_{nk} at a time. If we consider superpositions $\psi_k = \sum_n c_n * c_n \psi_{nk}$, we simply need to replace $K_{nk}{}^2$ with $\langle K_{nk} \rangle^2$. Using Eq. (2.3. 9) we can say $|\mathbf{p}_{nk}| = n\hbar k$ and $\langle |\mathbf{p}_k| \rangle = \langle n \rangle \hbar k \; = \hbar k \sum_n c_n * c_n \cdot n$, then using Eq. (3.1. 11)

$$\langle K_{nk} \rangle = \frac{\lambda\!\!\!^{-}_c k \sqrt{2s}}{2} \langle n \rangle \;\&\; \langle K_{nk} \rangle^2 = \frac{\lambda\!\!\!^{-}_c{}^2 k^2 s}{2} \langle n \rangle^2 \;\text{ where }\; \langle n \rangle = \sum_n (c_n * c_n) n \tag{3.2. 10}$$

Replacing $K_{nk}{}^2$ with $\langle K_{nk} \rangle^2 = \frac{\lambda\!\!\!^{-}_c{}^2 k^2 s}{2} \langle n \rangle^2$ in the key equations (3.2. 3) & (3.2. 6) we find there is no change to the final results. The laws of quantum mechanics tell us that the total angular momentum has to be precisely either integral $\hbar$ or half integral $\hbar/2$. If we look at the above integrals used to derive the total angular momentum we can see that N must be 1, also s must be exactly ½ or 1 for spin ½ & spin 1 massive particles respectively, in Eq. (2.2. 4) our probability formula. Also these integrals are infinite sums of positive and negative integral $\hbar$ that are virtual and cannot be observed, and occur with probabilities between 0 & 1. If an infinite superposition for an electron is in a spin up state and flips to spin down in a magnetic field, a real $m = \pm 1$ photon is emitted carrying away the change in angular momentum. This is the only real effect observed from this infinity of $(l = 3, m = +2)$ virtual wavefunctions all flipping to $(l = 3, m = -2)$ states, plus an infinite flipping of the virtual zero point vector debts. Also Eq's. (3.2. 3) & (3.2. 6) are true *only if our high energy cutoff is at infinity and the low frequency cutoff is at zero*. We will look at high frequency cutoffs in section 4.2 where we find they cutoff near the Planck scale and in section 6.1 low frequency cutoffs near the radius of the causally connected horizon depending on whether they are $N = 1$ or 2.

3.3 Ratios between Primary and Secondary Coupling

3.3.1 Initial simplifying assumptions

This section is based on a special case "thought experiment" to try and illustrate as simply as possible how superpositions interact with one another in the same way as virtual photons interact with electrons for example. It is also important to remember here that *because* *primary coupling constants are to bare charges (section 2.2.2), and thus fixed for all k, while secondary coupling constants run with k, that the coupling ratios can be defined only at the cutoff value of k applying to the bare charge (sections 4.1.1 & 4.3).* Initially it might seem there must be three ratios, χ_G for gravity, χ_C for colour and χ_{EM} for electromagnetism. However there is a very simple relationship between the colour and electromagnetic ratios and we will also find that $\chi_G = \chi_C$. We define these three ratios as follows:

$$\frac{\alpha_{Gravity(\text{Secondary})}}{\alpha_{Gravity(\text{Primary})}} = \frac{G_{\text{Sec}}}{G_{\text{Pri}}} = \frac{1}{\chi_G},\ \frac{\alpha_{Colour(\text{Secondary})}}{\alpha_{Colour(\text{Primary})}} = \frac{\alpha_{3S}}{\alpha_{3P}} = \frac{1}{\chi_C},\ \frac{\alpha_{EM(\text{Secondary})}}{\alpha_{EM(\text{Primary})}} = \frac{\alpha_{EMS}}{\alpha_{EMP}} = \frac{1}{\chi_{EM}}. \qquad (3.3.\ 1)$$

From Table 2.2. 1 there are 6 fundamental primary charges for electrons and positrons. But electrons and positrons are defined as fundamental charges. In other words what we define as a fundamental electric charge is in reality 6 primary charges. Of course we never measure 6 as their effect is reduced by the ratio between primary and secondary coupling. Because electromagnetic and colour coupling are both via spin one bosons, their coupling ratios are essentially the same but related simply as $6^2 = 36:1$.

$$\left[\frac{\alpha_{EM(\text{Secondary})}}{\alpha_{EM(\text{Primary})}} = \frac{\alpha_{EMS}}{\alpha_{EMP}} = \frac{1}{\chi_{EM}}\right] = 36 \times \left[\frac{\alpha_{Colour(\text{Secondary})}}{\alpha_{Colour(\text{Primary})}} = \frac{\alpha_3}{\alpha_{3P}} = \frac{1}{\chi_C}\right] \qquad (3.3.\ 2)$$

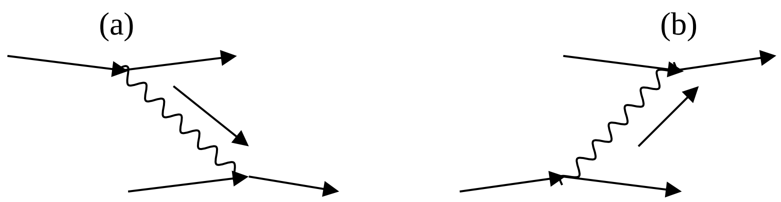
The secondary couplings α_{EMS} & α_{3S} are the bare charge values, and the gravitational coupling G_S is the usual measured value between Planck masses; all three are at the superposition cutoff near the Planck length. We are assuming the secondary gravitational coupling constant does not run with wavenumber *k* as the colour and electromagnetic do (see section 4.2.1). Also we assumed in section 2.2.2 that $\alpha_{3P} = 1;$ thus from Eq. (3.3. 2)

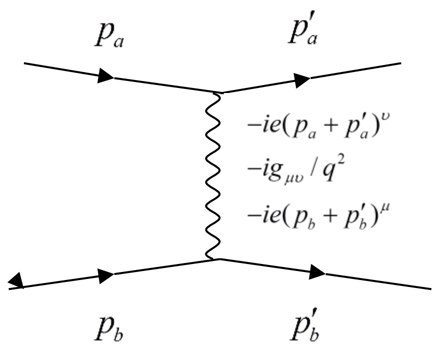
$$\alpha_{3S}^{-1} = \chi_C \qquad (3.3.\ 3)$$

In other words provided $\alpha_{3P} = 1,$ the ratio χ_C is also the inverse of the colour coupling constant α_3 at the superposition cutoff near the Planck length L_p (section 4.3).

From the above paragraphs to find the coupling ratios we need secondary interactions that are between bare charges. But this implies extremely close spacing where the effects of spins dominate. If the spacing is sufficiently large the effects of spin can be ignored but then we are not looking at bare charges. However we can ignore the effects of shielding due to virtual charged pairs by imagining as a simple thought experiment, an interaction between bare charges even at such large spacing. We can also simplify things further by considering only scalar or coulomb type interactions at this large spacing. We are also going to temporarily ignore Eq. (3.3. 2) and imagine that we have only one primary electric and or one colour charge. Consider two infinite superpositions and (*due to the above simplifying assumptions*) imagine them as spin zero charges. QED considers the interaction between them as a single covariant combination of two separate and opposite direction non-covariant interactions (a) plus (b) as in the Feynman diagram of Figure 3.3. 1 below. The Feynman transition amplitude is invariant in all frames [9]. So let us consider a special simple case in a CM frame where we have identical particles on a head on collision path with spatial momenta:

$$\mathbf{p}_a = -\mathbf{p}'_a = -\mathbf{p}_b = +\mathbf{p}'_b \tag{3.3. 4}$$





The Feynman diagram is drawn with a vertical photon line representing the superposition of two opposite direction and non covariant processes (a) plus (b).
The exchanged 4 momentum is:
$q = p_a - p'_a = p'_b - p_b$.

Figure 3.3. 1 Feynman diagram of virtual photon exchange between two spin zero particles of charge e.

From Eq. (3.3. 4) the initial and final spatial momenta are reversed with mirror images of each other at each vertex. Also in this simple special scalar case the transferred four momentum squared is simply the transferred three momentum squared.

$$q^2 = (p_a - p_a')^2 = (p_b - p_b')^2 = 4\mathbf{p}_a{}^2 = 4\mathbf{p}_b{}^2. \tag{3.3. 5}$$

If we look at Figure 3.3. 2 we see that at any fixed value of *k*, all modes ψ_{nk} in the groups of three overlapping superpositions for the various spins ½, 1 & 2 occupy very similar regions of space (provided they are all on the same centre.) The directions of their linear momenta are unknown but let us imagine some particular vector $\hbar\mathbf{k}$ that is parallel to the above vectors $\mathbf{p}_a = \mathbf{p}_b$. As we are considering only scalar interactions, all these modes must be spherically symmetric (as in section 3.2.2 for spins 1 & 2, and for spin ½ provided k or in turn β_{nk} is small enough the probability that it is not spherically symmetric can be extremely low) and at a fixed value of k they have momenta $\pm n\hbar\mathbf{k}$. Also as they overlap each other we can imagine units of $\pm\hbar\mathbf{k}$ quanta somehow transferring between these superpositions so that the values of n in each mode can change temporarily by ± 1 for times $\Delta T \approx \hbar / \Delta E$. The directions of these momentum transfers causing either repulsion or attraction depending on the charge signs of the superpositions at each vertex, whether the same or opposite.

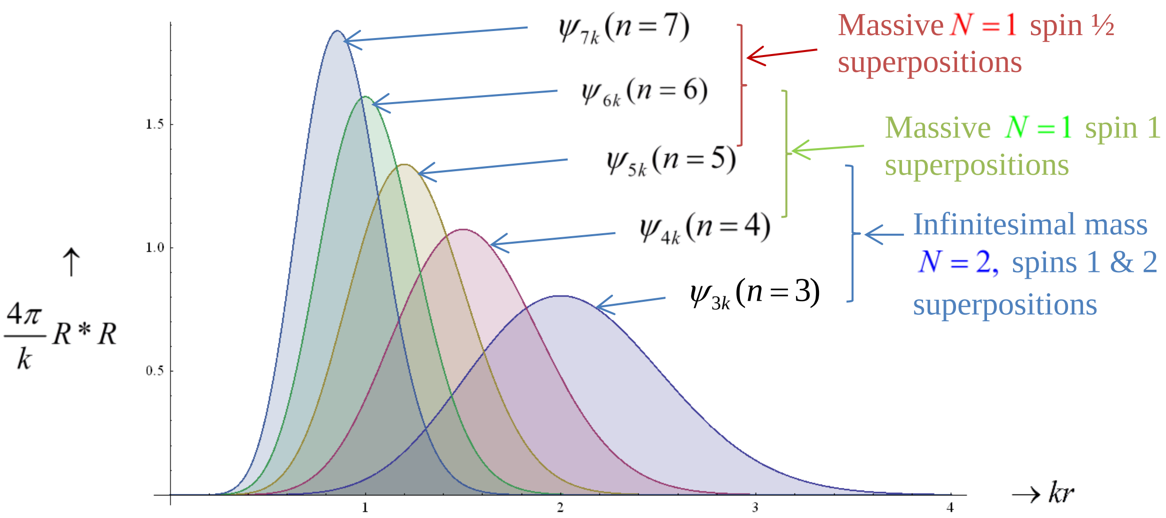


Figure 3.3. 2 All modes ψ_{nk} in the groups of three overlap at fixed wavenumber *k*.

3.3.2 Restrictions on possible Eigenvalue changes

Before we look at changing these Eigenvalues by $n = \pm 1$ we need to consider what restrictions there are on these changes.

From Eq. (2.3. 12) superposition ψ_k requires $Q^2A^2 = \sum_n c_n * c_n \frac{n^4\hbar^2k^4r^2}{81}$ and Eq.(2.2. 4) tells us the available $Q^2A^2 = \frac{\left[8+8\sqrt{\alpha_{EMP}}\right]^2}{3\pi sN}\hbar^2k^4r^2$ occurs with probability $= \frac{sN \cdot dk}{k}$.

For very brief periods the required value of Q^2A^2 can fluctuate, such as during these changes of momentum, but if its average value changes over the entire process then Eq. (2.2. 4) tells us that probability $sN \cdot dk / k$ changes also, and we have shown in section 3.2.1 that this is not allowed. For example in a spin ½ superposition $\psi_{5k}, \psi_{6k}, \psi_{7k}$, the average values of $|c_5|, |c_6|$ & $|c_7|$ must each remain constant. This can happen only if n remains within its pre-existing boundaries of $(5 \le n \le 7)$. For example if ψ_7 adds $+\hbar\mathbf{k}$, it can create ψ_8, but $|c_8|$ must average zero, which it can do only if it fluctuates either side of zero, and $|c_n|$ cannot be negative. Similarly $|c_4|$ must average zero, thus ψ_4 & ψ_8 are forbidden states. Keeping the average values of $|c_n|$ constant is also equivalent to a constant internal average particle energy (we have shown in section 3.2.1 that rest mass is a function of $\sum c_n * c_n . \mathbf{p}_{nk}{}^2$. By changing these Eigenvalues by $n = \pm 1$ there are only four possibilities; ψ_6 & ψ_7 can both reduce by $-\hbar\mathbf{k}$ quanta, ψ_6 & ψ_5 can both increase by $+\hbar\mathbf{k}$ quanta. If ψ_6 becomes ψ_7, $|c_7|$ also increases and $|c_6|$ decreases, but then ψ_7 has to drop back becoming ψ_6, with $|c_7|$ decreasing back down and $|c_6|$ increasing back up in exact balance. If we view this as one overall process the average values of both $|c_6|$ and $|c_7|$ remain constant but fluctuate continuously. We can use exactly the same argument if ψ_5 increases which has to be followed by ψ_6 dropping, where if we view this as one process again, the average values of both $|c_5|$ and $|c_6|$ remain constant. This is similar to a particle not being able to absorb a photon in a covariant manner, it has to reemit within time $\Delta T \approx \hbar / E$. With *spherical symmetry* the momentum $\mathbf{p} = \pm n\hbar\mathbf{k}$. If we change n by ± 1 the sign of $\mathbf{p} = \pm n\hbar\mathbf{k}$ determines the *direction of the momentum transfer* $\Delta\mathbf{p}$. In the above if $\psi_{5k} \to \psi_{6k}$ then returns $\psi_{6k} \to \psi_{5k}$, and $\mathbf{p} = \pm n\hbar\mathbf{k}$ *keeps the same sign during this process*, there is *no nett momentum transfer* and there is a probability of this, but it is not the probability we need. However if this process is as in figure Figure 3.3. 3.

$$\begin{array}{ccc} 6\hbar\mathbf{k} \rightarrow & & -5\hbar\mathbf{k} \leftarrow \\ \downarrow & \text{produces } \Delta\mathbf{p} = -\hbar\mathbf{k}\text{, but there is another } \Delta\mathbf{p} = -\hbar\mathbf{k} \text{ if returning via} & \downarrow \\ 5\hbar\mathbf{k} \rightarrow & & -6\hbar\mathbf{k} \leftarrow \\ |a\rangle & & |b\rangle \end{array}$$

Figure 3.3. 3

To get a net momentum transfer the momenta have to be in opposite directions for each half of this process. (Conservation of momentum allows this only if there is an equal and opposite transfer of momentum at the other vertex of the interaction.) The problem with this is that a total transfer of $\Delta\mathbf{p} = -2\hbar\mathbf{k}$ implies superpositions ψ_k interact with virtual $2k$ photons. Section 3.5 shows that interactions only with virtual k photons give the correct Dirac spin ½ magnetic energy. However just as transversely polarized photons are equal left and right polarization superpositions $|L\rangle/\sqrt{2}+|R\rangle/\sqrt{2}$, we can perhaps regard the Figure 3.3. 3 process as a similar equal superposition $|a\rangle/\sqrt{2}+|b\rangle\sqrt{2}$.

The figure 3.3.3 process becomes the superposition $\dfrac{|a\rangle}{\sqrt{2}}+\dfrac{|b\rangle}{\sqrt{2}}=\dfrac{|-\hbar\mathbf{k}\rangle}{\sqrt{2}}+\dfrac{|-\hbar\mathbf{k}\rangle}{\sqrt{2}}$ (3.3. 6)

We have two equal 50% probabilities of states $|a\rangle \& |b\rangle$ producing the required total $\Delta\mathbf{p} = -\hbar\mathbf{k}$. Also as from the above paragraphs the average values of $|c_5|$ and $|c_6|$ remain constant:

The probability of transitions $\psi_5 \rightleftarrows \psi_6$ must be the same in either direction. (3.3. 7)

As spherically symmetric states have momentum $\mathbf{p} = \pm n\hbar\mathbf{k}$:

We can also think of $\mathbf{p} = \pm n\hbar\mathbf{k}$ as a superposition $\mathbf{p} = |+n\hbar\mathbf{k}\rangle/\sqrt{2}+|-n\hbar\mathbf{k}\rangle/\sqrt{2}$. (3.3. 8)

3.3.3 Looking as just one vertex of the interaction first

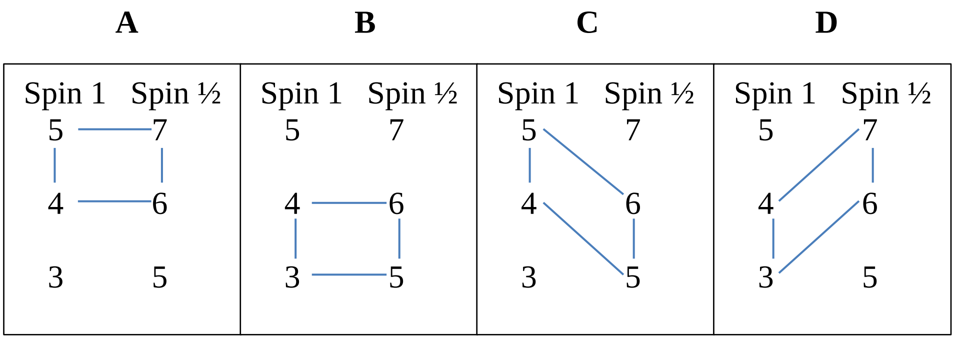
In Table 4.3. 1 and section 6.2 we introduce infinitesimal rest mass photons and gluons as superpositions of $\psi_{3k},\psi_{4k},\psi_{5k}$ where $N=2$ in Eq. (2.2. 4). Consider just one vertex of an infinitesimal rest mass spin 1 photon superposition $\psi_{3k},\psi_{4k},\psi_{5k}$ interacting with a spin ½ superposition $\psi_{5k},\psi_{6k,}\psi_{7k}$ at the same k. Looking at one possibility first, ψ_{4k} & ψ_{5k} for spin 1 and ψ_{6k} & ψ_{7k} for spin ½, we can apply the Figure 3.3. 3 process to get a nett momentum transfer. For this combination of Eigenfunctions there are four possible ways of getting the momentum transfer as in Figure 3.3. 4. In each of these 4 cases the amplitude for this to happen includes the factors $c_4 \cdot c_5 \cdot c_6 \cdot c_7$. Let us temporarily imagine $|c_4 \cdot c_5 \cdot c_6 \cdot c_7.| = 1$. Then $\mathbf{p} = +n\hbar\mathbf{k}$ as in $|a\rangle$ of Figure 3.3. 3 with an amplitude of $1/\sqrt{2}$ from Eq. (3.3. 8) transfers $\Delta\mathbf{p} = -\hbar\mathbf{k}$ also with an amplitude of $1/\sqrt{2}$, which is the required first half of our superposition Eq.(3.3. 6) $|a\rangle/\sqrt{2}+|b\rangle/\sqrt{2}$. Similarly $\mathbf{p} = -n\hbar\mathbf{k}$ as in $|b\rangle$ of Figure 3.3. 3 gives the second half. It would thus seem that our amplitude is simply $c_5 \cdot c_6 \cdot c_6 \cdot c_7$. However from Eq. (3.3. 7) there is a 50% probability of the transitions $\psi_5 \rightleftarrows \psi_6$ in either direction, or an extra $1/\sqrt{2}$ amplitude factor for $\psi_5 \rightleftarrows \psi_6$ in either direction, similarly an extra $1/\sqrt{2}$ amplitude factor for $\psi_6 \rightleftarrows \psi_7$. These two extra $1/\sqrt{2}$ factors reduce the amplitude $c_4 \cdot c_5 \cdot c_6 \cdot c_7$ to $c_4 \cdot c_5 \cdot c_6 \cdot c_7 / (\sqrt{2} \times \sqrt{2}) = c_4 \cdot c_5 \cdot c_6 \cdot c_7 / 2$. Thus adding the four cases in Figure 3.3. 4 together and *treating all other factors as 1*:

$$\text{Figure 3.3. 4 process amplitude factor is } 4 \times (c_4 \cdot c_5 \cdot c_6 \cdot c_7)/2 = 2c_4 \cdot c_5 \cdot c_6 \cdot c_7 \qquad (3.3.\ 9)$$

Spin 1 4 goes to 5 with **p** $\leftarrow$ 5 returns to 4 with **p** $\rightarrow$	Spin 1/2 7 goes to 6 with **p** $\leftarrow$ 6 returns to 7 with **p** $\rightarrow$	Spin 1 4 goes to 5 with **p** $\leftarrow$ 5 returns to 4 with **p** $\rightarrow$	Spin 1/2 6 goes to 7 with **p** $\rightarrow$ 7 returns to 6 with **p** $\leftarrow$
Spin 1 5 goes to 4 with **p** $\rightarrow$ 4 returns to 5 with **p** $\leftarrow$	Spin 1/2 7 goes to 6 with **p** $\leftarrow$ 6 returns to 7 with **p** $\rightarrow$	Spin 1 5 goes to 4 with **p** $\rightarrow$ 4 returns to 5 with **p** $\leftarrow$	Spin 1/2 6 goes to 7 with **p** $\rightarrow$ 7 returns to 6 with **p** $\leftarrow$

Figure 3.3. 4

The four possibilities in Figure 3.3. 4 are all between the same sets of Eigenfunctions ψ_{4k} & ψ_{5k} for spin 1, ψ_{6k} & ψ_{7k} for spin ½. But there are also four different sets of these between the groups of four Eigenfunctions A, B, C & D, as in Figure 3.3. 5; with their amplitudes from Eq. (3.3. 9) below each relevant box, which we also label as *A, B, C & D.* (Subscripts *a* refer to spin ½ and *b* to spin 1.)



Amplitudes: $A = 2c_{4b}c_{5b}c_{6a}c_{7a}$, $B = 2c_{3b}c_{4b}c_{6a}c_{5a}$, $C = 2c_{4b}c_{5b}c_{6a}c_{5a}$, $D = 2c_{3b}c_{4b}c_{6a}c_{7a}$.

Figure 3.3. 5

3.3.4 Assumptions when looking at both vertexes of the interaction

Because we are looking at an interaction between identical spin ½ fermions each vertex has the same groups of Eigenfunctions A,B,C&D as in Figure 3.3. 5. From section 2.2.2 and Figure 3.1. 4 the three Eigenfunctions forming each of the interacting particles are born simultaneously. It is thus reasonable to assume that the amplitudes of each group of three Eigenfunctions have the same complex phase angle. The two fermions and one boson can be at different complex phase angles to each other but each one individually is a superposition of three Eigenfunctions at the same complex phase angle. Thus the four amplitudes *A,B,C &D* from Figure 3.3. 5 (*A,B,C &D* each comprising two fermion amplitudes and two boson amplitudes) must all have the same complex phase angle. Similarly the four amplitudes $A', B', C' \& D'$ of vertex 2 in Figure 3.3. 6 all have a common complex phase angle.

Eigenfunction Groups	A	B	C	D
Vertex 1	Amplitude A	Amplitude B	Amplitude C	Amplitude D
Vertex 2	Amplitude A'	Amplitude B'	Amplitude C'	Amplitude D'

Figure 3.3. 6

We are also going to assume that Eigenfunctions A of vertex 1 interact only with Eigenfunctions A of vertex 2 and Eigenfunctions B of vertex 1 interact only with Eigenfunctions B of vertex 2 etc. Eigenfunctions A of vertex 1 do not interact with Eigenfunctions B of vertex 2 etc. Thus *if all other amplitude factors are 1*:

$$\text{The total interaction amplitude } = AA' + BB' + CC' + DD' \qquad (3.3.\ 10)$$

Apart from a different complex phase angle this is equivalent to: ($A \& A', B \& B'$ etc. all differ by the same complex phase angle.)

$$\text{Total interaction amplitude} = A^2 + B^2 + C^2 + D^2 \qquad (3.3.\ 11)$$

$$\text{Interaction probability} = (A^2 + B^2 + C^2 + D^2)*(A^2 + B^2 + C^2 + D^2) \qquad (3.3.\ 12)$$

Using $(A^2 * A^2) = (A*A)(A*A)$ etc. this is equivalent to

$$\text{Interaction probability} = (A^*A + B^*B + C^*C + D^*D)^2 \tag{3.3. 13}$$

Using $P_{5a} = c_{5a} * c_{5a}, P_{4b} = c_{4b} * c_{4b}$ etc. & $A^*A = 4P_{4b}P_{5b}P_{6a}P_{7a}$ etc. this is equivalent to

$$16\left[P_{4b}P_{5b}P_{6a}P_{7a} + P_{3b}P_{4b}P_{6a}P_{5a} + P_{4b}P_{5b}P_{6a}P_{5a} + P_{3b}P_{4b}P_{6a}P_{7a}\right]^2$$

$$= 16\left[P_{4b}(P_{3b} + P_{5b})\right]^2 \left[P_{6a}(P_{5a} + P_{7a})\right]^2$$

Then using $c_{3b} * c_{3b} + c_{4b} * c_{4b} + c_{5b} * c_{5b} = c_{5a} * c_{5a} + c_{6a} * c_{6a} + c_{6a} * c_{6a} = 1.$

$$\text{Interaction Probability} = 2^4\left[c_{4b} * c_{4b}(1 - c_{4b} * c_{4b})\right]^2 \left[c_{6a} * c_{6a}(1 - c_{6a} * c_{6a})\right]^2 \tag{3.3. 14}$$

We have assumed to here that *all other amplitude factors are 1*. However at each vertex there are both fermion and boson superposition probabilities from Eq. (2.2. 4). Writing the superposition probability at each vertex $sN \cdot dk / k$ as $s_{1/2}N_1 dk/k,\ s_1 N_2 dk/k$ for clarity (we will see in Table 4.3. 1 plus section 6.2 and its subsections that N can be either 1 or 2 for massive superpositions so we write: $\text{Spin } 1 = s_1,\ N = 1 \text{ is } N_1 \text{ etc.}$). Including these factors (*if all other factors are one*) in Eq. (3.3. 14) our overall probability at wavenumber k is

$$\left[\frac{2s_{1/2}N_1 c_{6a} * c_{6a}(1 - c_{6a} * c_{6a})}{k}\right]^2 \left[\frac{2s_1 N_2 c_{4b} * c_{4b}(1 - c_{4b} * c_{4b})}{k}\right]^2$$

$$= \frac{\left[2s_{1/2}N_1 c_{6a} * c_{6a}(1 - c_{6a} * c_{6a})\right]^2 \left[2s_1 N_2 c_{4b} * c_{4b}(1 - c_{4b} * c_{4b})\right]^2}{(k)^4}.$$

The momentum per transfer is a total of $\pm\hbar\mathbf{k}$ and using Eq's. (3.3. 5), (3.3. 6) & Figure 3.3. 3 we have $(\pm\hbar\mathbf{k})^4 = q^4$ (then putting $\hbar = 1$) the interaction probability:

$$= \frac{\left[2s_{1/2}N_1 c_{6a} * c_{6a}(1 - c_{6a} * c_{6a})\right]^2 \left[2s_1 N_2 c_{4b} * c_{4b}(1 - c_{4b} * c_{4b})\right]^2}{q^4} \tag{3.3. 15}$$

This is the scalar interaction probability between two spin ½ fermions exchanging infinitesimal rest mass spin 1 bosons at very large spacings, where the fermions are effectively spin zero, imagining them as bare charges, *with all other factors being one*. Going through exactly the same procedure but similarly exchanging spin 2 infinitesimal rest mass scalar gravitons (with $N = 2 = N_2$ for clarity) the *gravitational interaction probability* between fermions becomes (using *subscript c* for spin 2) *if all other amplitude factors are 1*:

$$= \frac{\left[2s_{1/2}N_1 c_{6a} * c_{6a}(1-c_{6a} * c_{6a})\right]^2 \left[2s_2 N_2 c_{4c} * c_{4c}(1-c_{4c} * c_{4c})\right]^2}{q^4} \text{ for fermions.} \tag{3.3. 16}$$

And if for example two spin 1 photons exchange spin 2 gravitons (all infinitesimal rest mass with $N = 2 = N_2$) the *interaction probability* becomes *if all other amplitude factors are 1*:

$$= \frac{\left[2s_1 N_2 c_{4b} * c_{4b}(1-c_{4b} * c_{4b})\right]^2 \left[2s_2 N_2 c_{4c} * c_{4c}(1-c_{4c} * c_{4c})\right]^2}{q^4} \text{ for } N = 2 \text{ photons.} \tag{3.3. 17}$$

If two massive $N = 1$ photons (as in Figure 3.3. 2) exchange spin 2 gravitons the *interaction probability* becomes *if all other factors are 1:*

$$= \frac{\left[2s_1 N_1 c_{5b} * c_{5b}(1-c_{5b} * c_{5b})\right]^2 \left[2s_2 N_2 c_{4c} * c_{4c}(1-c_{4c} * c_{4c})\right]^2}{q^4} \text{ for } N = 1 \text{ photons.} \tag{3.3. 18}$$

If two spin 2 gravitons exchange virtual gravitons (all infinitesimal rest mass with $N = 2$ and spin $s = 2$) *the interaction probability* becomes *if all other amplitude factors are 1*:

$$= \frac{\left[2s_2 N_2 c_{4c} * c_{4c}(1-c_{4c} * c_{4c})\right]^2 \left[2s_2 N_2 c_{4c} * c_{4c}(1-c_{4c} * c_{4c})\right]^2}{q^4} \text{ for } N = 2 \text{ gravitons.} \tag{3.3. 19}$$

General Relativity section 1.1.1 tells us the emission of gravitons is identical for both mass and energy. So keeping all other factors (such as mass) in Eq's. (3.3. 16), (3.3. 17) , (3.3. 18) & (3.3. 19) constant, the exchange probabilities must be the same in all four, implying:

$$2s_2N_2c_{4c}*c_{4c}(1-c_{4c}*c_{4c}) = 2s_1N_2c_{4b}*c_{4b}(1-c_{4b}*c_{4b}) = 2s_1N_1c_{5b}*c_{5b}(1-c_{5b}*c_{5b})$$
$$= 2s_{1/2}N_1c_{6a}*c_{6a}(1-c_{6a}*c_{6a}) \quad \text{or}$$

$$\underset{N=2\ \text{Spin } 2}{8c_{4b}*c_{4b}(1-c_{4b}*c_{4b})} = \underset{N=2\ \text{Spin } 1}{4c_{4b}*c_{4b}(1-c_{4b}*c_{4b})} = \underset{N=1\ \text{Spin } 1}{2c_{5b}*c_{5b}(1-c_{5b}*c_{5b})} = \underset{N=1\ \text{Spin } 1/2}{c_{6a}*c_{6a}(1-c_{6a}*c_{6a})} \qquad (3.3.\ 20)$$

Assuming all other factors are 1, and remembering that we are simplifying by looking at spin ½ superpositions sufficiently far apart so we can treat them as approximately spherically symmetric or effectively spin zero even though they are supposed to be bare charges:
Under the same scalar exchange conditions QED tells us that with electrons for example:

The probability of scalar or coulomb exchange in Eq (3.3. 15). $= \dfrac{4\alpha^2}{q^4}$. (3.3. 21)

Let us also temporarily ignore the fact that gluons have very limited range and imagine that our "thought experiment" also applies to colour charges exchanging gluons. Now treat *α as a general coupling constant regardless of type of charge or whether it is mass. Also temporarily ignore the fact that there can be up to six primary electric charges. Plus call the general coupling ratio in all three cases χ.* If the primary constant is one and putting $\alpha = \chi^{-1}$, $2s_{1/2} = 1$, $2s_1 = 2$, $N_1 = 1$ & $N_2 = 2$ we can equate Eq's. (3.3. 15) & (3.3. 21):

$$\frac{[c_{6a}*c_{6a}(1-c_{6a}*c_{6a})]^2[4c_{4b}*c_{4b}(1-c_{4b}*c_{4b})]^2}{q^4} = \frac{4(\chi^{-1})^2}{q^4} \qquad (3.3.\ 22)$$
$$\text{or} \quad [c_{6a}*c_{6a}(1-c_{6a}*c_{6a})][c_{4b}*c_{4b}(1-c_{4b}*c_{4b})] = \chi^{-1}/2$$

Combining Eq's. (3.3. 20) & (3.3. 22)

$$\overset{N=2\ \text{Spin } 2}{8c_{4c}*c_{4c}(1-c_{4c}*c_{4c})} = \overset{N=2\ \text{Spin } 1}{4c_{4b}*c_{4b}(1-c_{4b}*c_{4b})} = \overset{N=1\ \text{Spin } 1}{2c_{5b}*c_{5b}(1-c_{5b}*c_{5b})} = \overset{N=1\ \text{Spin } 1/2}{c_{6a}*c_{6a}(1-c_{6a}*c_{6a})}$$
$$= \sqrt{2/\chi} \qquad (3.3.\ 23)$$

The coupling ratio is thus fundamentally the same for colour, electromagnetism and gravity. However reintroducing six primary electric charges Eq. (3.3. 2) and because of the way electric charge is defined

$$\frac{1}{\chi}=\left[\frac{1}{\chi_G}=\frac{G_{\text{Secondary}}}{G_{\text{Primary}}}\right]=\left[\frac{1}{\chi_C}=\frac{\alpha_{Colour(\text{Secondary})}}{\alpha_{Colour(\text{Primary})}}\right]=\frac{1}{36}\left[\frac{1}{\chi_{EM}}=\frac{\alpha_{EM(\text{Secondary})}}{\alpha_{EM(\text{Primary})}}\right] \qquad (3.3.\ 24)$$

We also need to remember from Eq. (3.3. 3) that $\alpha_{3S}{}^{-1}=\chi_C$, or in other words provided $\alpha_{3P}=1$, the ratio χ_C is also the inverse of the colour coupling constant at the superposition cutoff near the Planck length (section 4.3).

Equations (3.3. 22) & (3.3. 23) tell us that for any interactions between pairs of superpositions, the inverse coupling ratio always involves the product of the central superposition member probability, by one minus that probability, and the same product for the other superposition. In section 4.3 when we introduce gravity we will solve these ratios and show that despite all the simplifications the above equations seem to fit the Standard Model reasonably well provided there are only three families of fermions.

3.4 Electrostatic Energy between two Infinite Superpositions

3.4.1 Using a simple quantum mechanics pre QED era approach

In section 3.3 we have shown that fermion superpositions can exchange boson superpositions in the same way as electrons can exchange virtual photons for example. Providing the superposition amplitudes are appropriate, the coupling constants can be just as in QED, though we will look further at this in section 4.1.1. So it would initially seem that evaluating electrostatic energy between superpositions is unnecessary. However there are three key reasons for revisiting a quantum mechanics method used before the days of QED to find the scalar potentials between two charges (or infinite superpositions).

a) It facilitates a simplified solution to the magnetic energy between superpositions in section 3.5 where we modify relevant equations in a simple manner.
b) We also use some of the same equations when looking at why borrowing energy and mass from zero point fields requires the universe to expand after the Big Bang.
c) These same equations also help explain in a very simple manner why General Relativity requires the metric to change near any mass concentration.

We assume spherically symmetric $l=3$ superpositions emit virtual scalar photons in this section and $l=3, m=\pm 2$ superpositions emit virtual $m=\pm 1$ photons in section 3.5. As section 3.3 has shown that we can achieve the same electromagnetic coupling constant α we can use the scalar photon emission probability $(2\alpha/\pi)(dk/k)$ covered in section 2.1.1. From section 3.3 we can also see that the effective average emission point has to be the center of superpositions. The probability of finding this interacting virtual photon (or spin 1 superposition) decays exponentially with distance travelled. The normalized wavefunction ψ for such a virtual scalar photon of wave number k emitted at $r=0$ is:

$$\psi = \sqrt{\frac{2k}{4\pi}}\frac{e^{-kr}e^{+i\omega t}}{r} = \sqrt{\frac{2k}{4\pi}}\frac{e^{-kr}e^{+ikr}}{r}.$$

Wavefunction ψ is spherically symmetric as scalar photons are time polarized**.** Figure 3.4. 1 plots the radial probabilities of the exponential range of the virtual photon and the dominant $n=6$ mode of its relating superposition ψ_k. The effective range of the interacting photon is of a similar order to the radial probability dimensions of ψ_{6k}.

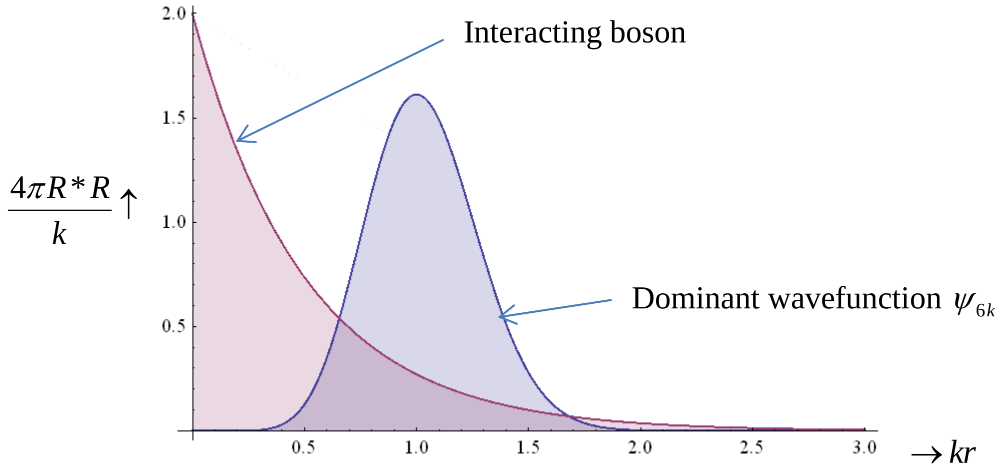


Figure 3.4. 1 Radial probability of ψ_{6k} and the exponential decay with radius of its interacting virtual boson $R^*R \propto 2ke^{-2kr}$. These curves are the same for all k, applying equally to virtual photons, gravitons and to large k value gluons etc.

For simplicity in what follows we locate two superpositions (which we refer to as sources) in cavities that are small in relation to the distance between them. The accuracy of our results depends on how far apart they are in relation to the cavity size. Consider two spherically symmetric sources distance $2C$ apart emitting virtual scalar photons as in Figure 3.4. 2 where point P is r_1 from source 1, & r_2 from source 2. Let ψ_1 be the amplitude from source 1, and ψ_2 be the amplitude from source 2.

Thus $\quad \psi_1 = \sqrt{\dfrac{2k}{4\pi}}\dfrac{e^{-kr_1+ikr_1}}{r_1} \quad \& \quad \psi_2 = \sqrt{\dfrac{2k}{4\pi}}\dfrac{e^{-kr_2+ikr_2}}{r_2}$ (3.4. 1)

Consider $\quad (\psi_1+\psi_2)*(\psi_1+\psi_2) \quad = \psi_1*\psi_1+\psi_1*\psi_2+\psi_2*\psi_1+\psi_2*\psi_2$

Now $\psi_1*\psi_1$ & $\psi_2*\psi_2$ are just the normal probability densities around sources 1 & 2 as though they are infinitely far apart but the work done per pair of superpositions k on bringing 2 sources closer together is in the interaction term: $\psi_1*\psi_2+\psi_2*\psi_1$.

$$\psi_1*\psi_2 = \frac{2k}{4\pi r_1 r_2}e^{-kr_1}e^{-kr_2}e^{+ikr_1}e^{-ikr_1} = \frac{2k}{4\pi r_1 r_2}e^{-k(r_1+r_2)}e^{+ik(r_1-r_2)}$$

$$\psi_2*\psi_1 = \frac{2k}{4\pi r_1 r_2}e^{-kr_2}e^{-kr_1}e^{+ikr_2}e^{-ikr_2} = \frac{2k}{4\pi r_1 r_2}e^{-k(r_1+r_2)}e^{-ik(r_1-r_2)}$$

$$\psi_1*\psi_2+\psi_2*\psi_1 = \frac{2k}{4\pi r_1 r_2}e^{-k(r_1+r_2)}\left[e^{+ik(r_1-r_2)}+e^{-ik(r_1-r_2)}\right]$$

$$= \frac{4k}{4\pi r_1 r_2}e^{-k(r_1+r_2)}\cos k(r_1-r_2)$$

Now put $(A=r_1+r_2, B=r_1-r_2)$ & $\quad \psi_1*\psi_2+\psi_2*\psi_1 = \dfrac{4k}{4\pi r_1 r_2}e^{-Ak}\cos kB$ (3.4. 2)

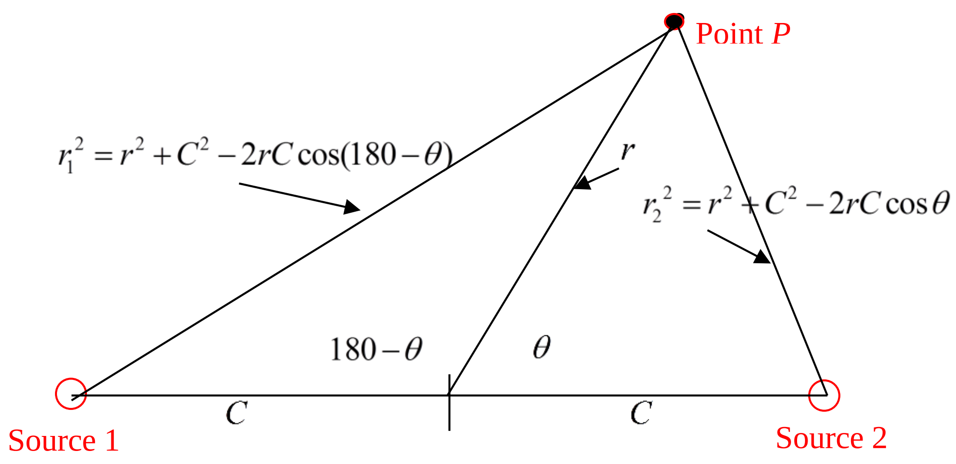


Figure 3.4. 2

Real work is done bringing superpositions together and we can treat these virtual photons as having real energy $\hbar\omega = \hbar kc$. Using virtual photon emission probability $(2\alpha/\pi)(dk/k)$ from section 2.1.1.

$$\text{Energy } \hbar\omega = \hbar kc \times \text{Probability } \frac{2\alpha}{\pi}\frac{dk}{k} = \frac{2\alpha\hbar c}{\pi}dk \tag{3.4. 3}$$

Our interaction energy @ k is thus $(\psi_1 {}^*\psi_2 + \psi_2 {}^*\psi_1\ )\dfrac{2\alpha\hbar c}{\pi}dk$ which using Eq. (3.4. 2)

Interaction energy @ k is $\dfrac{2\alpha\hbar c}{\pi}dk \ \ \dfrac{4k}{4\pi r_1 r_2}e^{-Ak}\cos kB$. The total interaction energy density due to $\psi_1 {}^*\psi_2 + \psi_2 {}^*\psi_1$ for all k is

$$\frac{2\alpha\hbar c}{\pi}\frac{4}{4\pi r_1 r_2}\int_0^\infty ke^{-Ak}\cos(Bk)dk \tag{3.4. 4}$$

$$\int_0^\infty ke^{-Ak}\cos(Bk)dk = \frac{A^2 - B^2}{(A^2 + B^2)^2} \tag{3.4. 5}$$

Where $\quad A^2 = (r_1 + r_2)^2 = r_1^2 + 2r_1 r_2 + r_2^2 \ \ \& \ \ B^2 = (r_1 - r_2)^2 = r_1^2 - 2r_1 r_2 + r_2^2$

$$\text{Thus} \quad A^2 = (r_1 + r_2)^2 = r_1^2 + 2r_1 r_2 + r_2^2 \ \ \& \ A^2 + B^2 = 2(r_1^2 + r_2^2) \tag{3.4. 6}$$

$$= 2(r^2 + C^2) \text{ as } \cos(180 - \theta) = -\cos\theta$$

$$\text{and} \quad A^2 + B^2 = 4(r^2 + C^2) \tag{3.4. 7}$$

Putting Eq's. (3.4. 4), (3.4. 5), (3.4. 6) & (3.4. 7) together $\dfrac{A^2 - B^2}{(A^2 + B^2)^2} = \dfrac{4r_1 r_2}{16(r^2 + C^2)^2}$

$$\int_0^\infty ke^{-Ak}\cos(Bk)dk = \frac{r_1 r_2}{4(r^2 + C^2)^2}$$

$$\frac{2\alpha\hbar c}{\pi}\frac{4}{4\pi r_1 r_2}\int_0^\infty ke^{-Ak}\cos(Bk)dk = \frac{2\alpha\hbar c}{\pi}\frac{4}{4\pi r_1 r_2}\frac{r_1 r_2}{4(r^2 + C^2)^2}$$

$$= \frac{2\alpha\hbar c}{\pi}\frac{1}{4\pi}\frac{1}{(r^2 + C^2)^2} \tag{3.4. 8}$$

This is the total interaction energy density of time polarized virtual photons at point P due to $\psi_1 * \psi_2 + \psi_2 * \psi_1$ for all k and there are no directional vectors to take into account. We will use similar equations for the vector potential ($m = \pm 1$) photons for magnetic energies but will then need directional vectors. Equation (3.4. 8) is the energy due to the interaction of amplitudes at any radius r from the centre of the pair. It is independent of θ, and to get the total energy of interaction we multiply by $4\pi r^2 dr$ for any layer dr and integrate from $r = 0 \to \infty$.

The total interaction energy is $\displaystyle \frac{2\alpha\hbar c}{\pi}\int_0^\infty \int_0^\infty (\psi_1 * \psi_2 + \psi_2 * \psi_1) dk \; 4\pi r^2 dr$

Using Eq. (3.4. 8) $\displaystyle = \frac{2\alpha\hbar c}{\pi}\frac{1}{4\pi}\int_0^\infty \frac{4\pi r^2 dr}{(r^2 + C^2)^2}$

Thus $\displaystyle \frac{2\alpha\hbar c}{\pi}\int_0^\infty \int_0^\infty (\psi_1 * \psi_2 + \psi_2 * \psi_1) dkdv = \frac{2\alpha\hbar c}{\pi}\int_0^\infty \frac{r^2 dr}{(r^2 + C^2)^2}$

$$\int_0^\infty \frac{r^2 dr}{(r^2 + C^2)^2} = \frac{1}{2C}\frac{\pi}{2}$$

The interaction or potential energy is $\displaystyle \frac{\alpha\hbar c}{2C} = \frac{\alpha\hbar c}{R}$ (3.4. 9)

If $R = 2C$ is the distance between the centres of our assemblies, this is the classical potential. The procedure used here with small changes, simplifies the derivation of the magnetic moment; we reuse some equations, but in a slightly modified form taking polarization vectors into account.

3.5 Magnetic Energy between two spin aligned Infinite Superpositions

In this section we are going to consider two infinite superpositions that form Dirac spin ½ states. We will look at the magnetic energy between them when they are both in a spin up state say along some z axis as in Figure 3.5. 1. *We are not looking at the magnetic energy here when they are both coupled in a spin 0 or spin 1 state*. That is, both Dirac spin ½ states have their $\sqrt{3}\hbar/2$ spin vectors randomly oriented around the z axis with $\hbar/2$ components aligned along this z axis. Also in this section we will be dealing with transversely polarized virtual photons and must take account of polarization vectors. In section 3.2.2 and Eq. (3.2. 7) spin ½ states are generated only from $l=3, m=2$ states and as transversely polarized photons are superpositions of $m=\pm 1$ photons they can only be emitted from these $l=3, m=2$ states, the remaining states are spherically symmetric and cannot emit transversely polarized photons. We don't yet know the value of amplitudes $|c_n|$ so we will derive the magnetic energy in terms of these. We will then equate this energy to the Dirac values assuming a g value of 2 before QED corrections; this allows us to evaluate in section 4.4 the amplitudes $|c_n|$ in terms of the ratio χ_{EM} between primary and secondary electromagnetic coupling. We can then evaluate in section 4.1 the primary electromagnetic coupling constant α_{EMP} in terms of the ratio χ_{EM}. This section uses the same format as Chapter 18, "The Feynman Lectures on Physics" Volume 3, Quantum Mechanics [11].

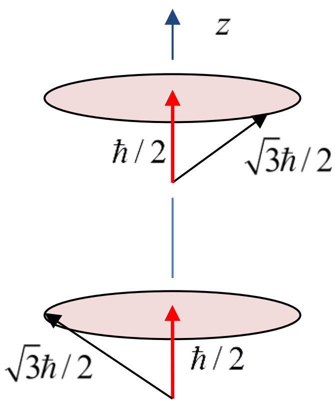


Figure 3.5. 1

An $l=3, m=2$ state can emit a right hand circularly (R.H.C.) polarized $(m=+1)$ photon in the $+z$ direction. Let the amplitude for this be temporarily $|R\rangle$.
An $l=3, m=-2$ state can emit a left hand circularly (L.H.C.) polarized $(m=-1)$ photon in the $+z$ direction. Let the amplitude for this also be temporarily $|L\rangle$.
First rotate the z axis about the y axis by angle θ (call this operation $S|R\rangle$) then use $\langle x'| = (1/\sqrt{2})\left[\langle R'| + \langle L'|\right]$ and multiply on the right by operation $S|R\rangle$.
The amplitude to emit a transversely polarized photon in the x' direction is thus

$$\langle x'|S|R\rangle = \frac{1}{\sqrt{2}}\left[\langle R'|S|R\rangle + \langle L'|S|R\rangle\right]$$

Where $\langle R'|S|R\rangle = \langle 3,+2'|S|3,+2\rangle = (1/4)\left[2+2\cos\theta - 4\sin^2\theta + 3\sin^2\theta\cos\theta\right]$ is the amplitude an $l=3, m=2$ state remains in an $l=3, m=2$ state after rotation by angle θ.

Also $\langle L'|S|R\rangle = -\langle 3,-2'|S|3,+2\rangle = (1/4)\left[2-2\cos\theta - 4\sin^2\theta - 3\sin^2\theta\cos\theta\right]$ is minus the amplitude that an $l=3, m=2$ state is in an $l=3, m=-2$ state after rotation by θ.

Putting this together
$$\langle x'|S|R\rangle = \frac{1-2\sin^2\theta}{\sqrt{2}} = \frac{\cos 2\theta}{\sqrt{2}} \tag{3.5. 1}$$

An $l=3, m=2$ state can also emit an $(m=+1)$ photon in the $-z$ direction but it will now be left hand circularly polarized. Let this amplitude be temporarily: $|L\rangle$.
Similarly an $l=3, m=-2$ state can emit an $(m=-1)$ photon in the $-z$ direction which is right hand circularly polarized. Let this amplitude be temporarily: $|R\rangle$.

We can go through the same procedure as above to get $\langle x'|S|L\rangle = \dfrac{\cos 2\theta}{\sqrt{2}}$ (3.5. 2)

This amplitude Eq. (3.5. 2) is for a photon emitted in the opposite direction to amplitude Eq. (3.5. 1) but $\cos 2\theta = \cos 2(180+\theta)$ and we can simply add these two amplitudes. Let us assume however that an $l=3, m=2$ state has equal amplitudes to emit in the $+z$ & $-z$ directions of $|R\rangle/\sqrt{2}$ and $|L\rangle/\sqrt{2}$.

With these amplitudes;
$$\frac{1}{\sqrt{2}}\left[\langle x'|S|R\rangle + \langle x'|S|L\rangle\right] = \frac{\cos 2\theta}{2} + \frac{\cos 2\theta}{2} = \cos 2\theta \tag{3.5. 3}$$

Eqation (3.5. 3) is the angular component of the amplitude for a transverse x' polarization in the new z' direction where $x \to x'$ & $z \to z' = \theta$. When $\theta = 0$ or 180 the on axis amplitude for transverse polarization is one as expected ignoring other factors. Using the same normalization factors (we check the validity of this in section 3.5.1 we can still use the amplitudes and phasing of our original time mode photons Eq's. (3.4. 1) but instead of including polarization vectors we will for simplicity just use the cosine of the angle $(\gamma-\delta)$ between them (as in Figure 3.5. 2) as a multiplying factor. Including the angular factor Eq. (3.5. 3) in our earlier scalar amplitudes Eq's. (3.4. 1) we have for our new wavefunctions:

$$\psi_1 = \cos 2\delta\sqrt{\frac{2k}{4\pi}}\frac{e^{-kr_1+ikr_1}}{r_1} \quad \& \quad \psi_2 = \cos 2\gamma\sqrt{\frac{2k}{4\pi}}\frac{e^{-kr_2+ikr_2}}{r_2} \tag{3.5. 4}$$

The transverse polarized photons from sources (1) & (2) have polarization vectors $|x_1\rangle$ and $|x_2\rangle$ at angle to each other $(\gamma-\delta)$, (Figure 3.5. 2) and the complex product becomes:

$$(\psi_1+\psi_2)^*(\psi_1+\psi_2) \quad =\psi_1{}^*\psi_1+(\psi_1{}^*\psi_2+\psi_2{}^*\psi_1)(\cos(\gamma-\delta)+\psi_2{}^*\psi_2$$

Where the interaction term is now: $(\psi_1{}^*\psi_2+\psi_2{}^*\psi_1)\cos(\gamma-\delta)$ and as in the scalar case (section 3.4.1) but now using Eq's. (3.5. 4)

$$\psi_1{}^*\psi_2\cos(\gamma-\delta)=\cos 2\delta\cos 2\gamma\frac{2k}{4\pi r_1 r_2}e^{-k(r_1+r_2)}e^{+ik(r_1-r_2)}\cos(\gamma-\delta)$$

$$\psi_2{}^*\psi_1\cos(\gamma-\delta)=\cos 2\delta\cos 2\gamma\frac{2k}{4\pi r_1 r_2}e^{-k(r_1+r_2)}e^{-ik(r_1-r_2)}\cos(\gamma-\delta)$$

$$(\psi_1{}^*\psi_2+\psi_2{}^*\psi_1)\cos(\gamma-\delta)=\cos 2\delta\cos 2\gamma\frac{4k}{4\pi r_1 r_2}e^{-Ak}\cos kB\,\cos(\gamma-\delta) \qquad (3.5.\ 5)$$

Where just the same as section 3.4.1, Eq. (3.4. 2) we have used: $A=r_1+r_2\ \&\ B=r_1-r_2$.
In the laboratory frame ψ_{nk} has amplitude β_{nk} to be in an $m=+2$ state (section 3.1). For a superposition ψ_k we need the expectation value $\langle\beta_{nk}\rangle$.

Using Eq's. (3.1. 11) & (3.1. 12) $\beta_{nk}=\dfrac{K_{nk}}{\sqrt{1+K_{nk}{}^2}}$ and $\langle\beta_{nk}\rangle=\left\langle\dfrac{K_{nk}}{\sqrt{1+K_{nk}{}^2}}\right\rangle$.

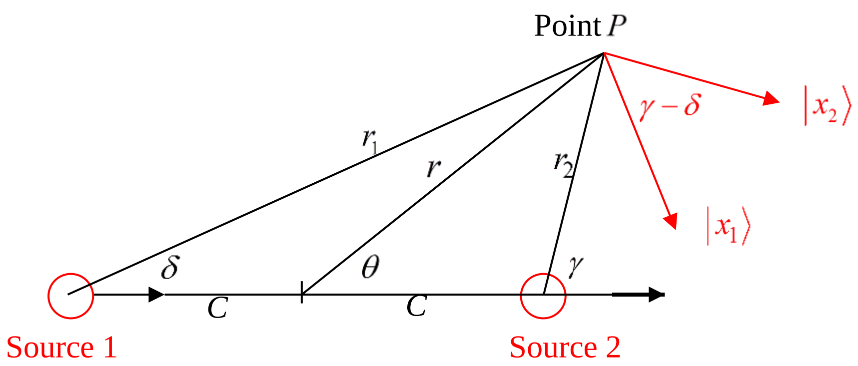


Figure 3.5. 2 Two sources $2C$ apart, both with $\beta_{nk}{}^2\times(m=+2)$ states along the joining line, $\delta\ \&\ \gamma$ are the respective angles to P, $r_1\ \&\ r_2$ are the respective distances to point P.

To keep our integrals simple we will assume that $\langle\beta_{nk}\rangle<<<<1$ or that our spacing is very large and that our interacting k values are also very small. (We can make a comparison with the Dirac values at any large spacing so our accuracy should not be affected.) Thus when $\left\langle \beta_{nk}{}^2 \right\rangle <<<< 1$ (and $\left\langle \gamma_{nk}{}^2 \right\rangle \approx 1$), we can approximate the above as: $\left\langle \beta_{nk} \right\rangle \approx \left\langle K_{nk} \right\rangle$ and from Eq. (3.2. 10) $\left\langle K_{nk} \right\rangle = \dfrac{\lambdabar_c k\sqrt{2s}}{2}\left\langle n \right\rangle \; = \dfrac{\lambdabar_c k}{2}\left\langle n \right\rangle$ as $s =$ ½ for spin ½ and also $\left\langle n \right\rangle = \sum_{n=5,6,7} (c_n{}^* c_n) n$

.

The approximate expectation value for the superposition is thus

$$\left\langle \beta_{nk} \right\rangle \approx \frac{\lambdabar_c \left\langle n \right\rangle k}{2} \tag{3.5. 6}$$

We will use the same probability of virtual photon emission and energy $\hbar kc$ as in Eq.(3.4. 3) of the scalar case.

$$\frac{2\alpha\hbar c}{\pi} dk \tag{3.5. 7}$$

3.5.1 Checking our normalization factors

Let us pause and check the reasonableness of this and our normalization factors. From Eq's. (3.4. 1) for scalar photons $\psi^*\psi = \dfrac{2k}{4\pi}\dfrac{e^{-2kr}}{r^2}$ × emission probability $\dfrac{2\alpha}{\pi}\dfrac{dk}{k}$ gives a scalar ψ_k emission probability density $\psi^*\psi \dfrac{4\alpha}{\pi} = \dfrac{2k}{4\pi}\dfrac{e^{-2kr}}{r^2}\dfrac{2\alpha}{\pi}\dfrac{dk}{k}$. The probability density at k for the transversely polarized case, using Eq. (3.5. 4) then including Eq. (3.5. 7) and $\left\langle \beta_{nk}{}^2 \right\rangle$

$$\left\langle \beta_{nk} \right\rangle^2 \psi'^*\psi' \frac{2\alpha}{\pi}\frac{dk}{k} = \left\langle \beta_{nk} \right\rangle^2 \cos^2 2\delta \frac{2k}{4\pi}\frac{e^{-2kr}}{r^2}\frac{2\alpha}{\pi}\frac{dk}{k} \; .$$

If we now consider the on axis $\delta = 0$ case we see that the transverse polarized on axis emission probability density at k is:

$$\left\langle \beta_{nk} \right\rangle^2 \frac{2k}{4\pi}\frac{e^{-2kr}}{r^2}\frac{2\alpha}{\pi}\frac{dk}{k} = \left\langle \beta_{nk} \right\rangle^2 \; \psi^*\psi \frac{2\alpha}{\pi}\frac{dk}{k}$$

Just as in QED the factor $\left\langle \beta_{nk} \right\rangle^2$ is the factor we need for this on axis emission probability density ratio between transverse and scalar polarization. This justifies using the same normalization constant $\sqrt{2k/4\pi}$ for both the scalar and magnetic wavefunctions. We seem to be on the right track and putting Eq. (3.5. 6) and Eq. (3.5. 7) into Eq. (3.5. 5) we get the transverse interaction energy @ wavenumber k :

$$\langle \beta_{nk} \rangle^2 (\psi_1{}^*\psi_2 + \psi_2{}^*\psi_1)\cos(\gamma-\delta)\left[\frac{2\alpha\hbar c}{\pi}dk\right]$$

$$=\left[\frac{\lambda\!\!\!^{-}{}_c{}^2 \langle n \rangle^2 k^2}{4}\right]\cos 2\delta \cos 2\gamma \frac{4k}{4\pi r_1 r_2} e^{-Ak} \cos kB \cos(\gamma-\delta)\left[\frac{2\alpha\hbar c}{\pi}dk\right]$$

Rearranging this: $\langle \beta_{nk} \rangle^2 (\psi_1{}^*\psi_2 + \psi_2{}^*\psi_1)\cos(\gamma-\delta)\left[\frac{2\alpha\hbar c}{\pi}dk\right]$

$$= \frac{2\langle n \rangle^2 \lambda\!\!\!^{-}{}_c{}^2 \alpha\hbar c}{\pi} \frac{\cos 2\delta \cos 2\gamma \cos(\gamma-\delta)}{4\pi r_1 r_2}\left[k^3 e^{-Ak}\cos(kB)dk\right] \tag{3.5. 8}$$

As in the scalar case we integrate over k first but now with a k^3 term due to the inclusion of the $\langle \beta_{nk} \rangle^2$ factor which is approximately proportional to k^2 from Eq. (3.5. 6).

Using $A = r_1 + r_2$ & $B = r_1 - r_2$ and Eq's. (3.4. 6) & (3.4. 7)

$$\int_0^\infty \left[k^3 e^{-Ak}\cos(kB)dk\right] = \frac{3}{8}\left[\frac{2r_1^2 r_2^2 - (r^2+C^2)^2}{(r^2+C^2)^4}\right]$$

And thus: $\displaystyle\int_0^\infty \langle \beta_{nk} \rangle^2 (\psi_1{}^*\psi_2 + \psi_2{}^*\psi_1)\cos(\gamma-\delta)\,\frac{2\alpha\hbar c}{\pi}dk$

$$= \frac{2\langle n \rangle^2 \lambda\!\!\!^{-}{}_c{}^2 \alpha\hbar c}{\pi} \frac{\cos 2\delta \cos 2\gamma \cos(\gamma-\delta)}{4\pi r_1 r_2} \times \frac{3}{8}\left[\frac{2r_1^2 r_2^2 - (r^2+C^2)^2}{(r^2+C^2)^4}\right] \tag{3.5. 9}$$

Equation (3.5. 9) is the magnetic interaction energy density at point P for all wave numbers k. Figure 3.5. 2 is a plane of symmetry that can be rotated through angle 2π around the axis of symmetry (the joining line along the axis of the 2 spin aligned sources). To evaluate the total magnetic energy density over all space we just multiply by $4\pi r^2 \sin\theta d\theta dr$. We thus integrate Eq. (3.5. 9) $\times\, 4\pi r^2 \sin\theta d\theta dr$ =

$$\frac{3\langle n \rangle^2 \lambda\!\!\!^{-}{}_c{}^2 \alpha\hbar c}{4\pi}\int_0^\infty\int_0^{\pi/2} \frac{\cos 2\delta \cos 2\gamma \cos(\gamma-\delta)}{r_1 r_2}\left[\frac{2r_1^2 r_2^2 - (r^2+C^2)^2}{(r^2+C^2)^4}\right] r^2 \sin\theta d\theta dr \tag{3.5. 10}$$

Now $\displaystyle\int_0^\infty\int_0^{\frac{\pi}{2}} \frac{\cos 2\delta \cos 2\gamma \cos(\gamma-\delta)}{r_1 r_2}\left[\frac{2r_1^2 r_2^2 - (r^2+C^2)^2}{(r^2+C^2)^4}\right] r^2 \sin\theta d\theta dr$ can be reduced to the single integral: $\frac{1}{8C^3}\int_0^1 \sqrt{1-x^2}\left[\frac{(7-5x^2)}{x^3}\ln\frac{1+x}{1-x}-\frac{14}{x^2}+\frac{16}{3}\right]dx$ which can be also expressed as an infinite series in p (to not confuse with superposition value n):

$$\frac{1}{8C^3}\sum_{p=1}^{p=\infty}\left[\frac{14}{2p+3}-\frac{10}{2p+1}\right]\frac{(2p-1)!}{(p-1)!(p+1)!4^p}\cdot\frac{\pi}{2} \quad =\frac{1}{8C^3}\frac{(160-51\pi)}{6}\cdot\frac{\pi}{2}$$

$$\text{(Putting } R=2C\text{)} \quad =\frac{1}{R^3}\frac{(160-51\pi)}{6}\cdot\frac{\pi}{2} \tag{3.5. 11}$$

$$\text{This infinite series is approximately } \approx -\frac{1}{R^3}\frac{\pi}{54(1.0045062....)} \tag{3.5. 12}$$

Putting Eq. (3.5. 12) into Eq. (3.5. 9) the total magnetic interaction energy over all frequencies and all space for 2 spin aligned infinite superpositions is:

$$U' \approx \frac{3\langle n\rangle^2 \lambda\!\!\!^{-}{}_c{}^2\alpha\hbar c}{4\pi}\left[-\frac{1}{R^3}\frac{\pi}{54(1.0045062....)}\right]$$

$$\text{We will call this} \quad U(\text{superpositions}) \approx -\left[\frac{\langle n\rangle^2 \lambda\!\!\!^{-}{}_c{}^2\alpha\hbar c}{72R^3(1.0045062....)}\right] \tag{3.5. 13}$$

We can equate this magnetic energy to the classical value assuming the Dirac value of $g=2$ for spin ½ (No QED corrections have been applied so it must be $g=2$). For the arrangement of spins as in Figure 3.5. 1 the Dirac magnetic energy between two spin ½ states is

$$U(\text{Dirac}) = -\left[\frac{2\mu^2}{4\pi\varepsilon_o c^2 R^3}\right] \tag{3.5. 14}$$

Using the Dirac magnetic moment $\mu=\frac{e\hbar}{2m_0}=\frac{e\hbar c}{2m_0 c}=\frac{ec\lambda\!\!\!^{-}{}_c}{2}$ the Dirac magnetic energy is

$$U(\text{Dirac}) = -\left[\frac{\lambda\!\!\!^{-}{}_c{}^2\alpha\hbar c}{2R^3}\right]$$

The approximation used in deriving Eq. (3.5. 6) $\gamma^2\beta^2 \approx \beta^2$ for $\beta^2 <<< 1$ is true only when $R >>> \lambda\!\!\!^{-}{}_c$. This error in β^2 is of the order of $\lambda\!\!\!^{-}{}_c{}^2/R^2$ and rapidly tends to zero with increasing R. There is no upper limit on the value of distance R we can choose. Thus comparing our estimate of the magnetic energy with Dirac's value when $R >>> \lambda\!\!\!^{-}{}_c$.

$$U(\text{Dirac}) = U(\text{Superpositions}) \text{ or } -\left[\frac{\lambda\!\!\!^{-}{}_c{}^2\alpha\hbar c}{2R^3}\right] \approx -\left[\frac{\langle n\rangle^2 \lambda\!\!\!^{-}{}_c{}^2\alpha\hbar c}{72R^3(1.0045062....)}\right] \qquad (3.5.\ 15)$$

All symbols cancel except $\langle n\rangle$ leaving: $\langle n\rangle^2 \approx 36(1.0045062.....)$

The expectation value $\langle n\rangle$ in our superposition is slightly more than $n=6$ our dominant mode. This is why we have used a three member superposition centred on this dominant $n=6$ mode. The two side modes $n=5$ & $n=7$ are smaller so that:

$$\langle n\rangle = \sum_{n=5,6,7.} (c_n c_n{}^*)n \approx \sqrt{36(1.0045062...)} \approx 6.01350345 \qquad (3.5.\ 16)$$

This is for Dirac spin ½ particles. This mean value of n creates a $g=2$ fermion which *QED corrections* (*which are secondary interactions*) *increase slightly to the experimental value.* In section 4.1 we will solve the primary electromagnetic coupling constant in terms of the ratio χ_{EM} using Eq. (3.5. 16). It is important to remember that this derivation of the magnetic energy applies only to two infinite assemblies (or particles) localized in two cavities that are small in relation to their distance R apart. They must also be both on the z axis with their spins aligned (or anti aligned) along this same z axis as in Figure 3.5. 1 & Figure 3.5. 2. Also the agreement with Dirac and in what follows is possible *if superposition* ψ_k *interacts only with virtual photons of the same wavenumber* k.

4 High Energy Superposition Cutoffs

4.1 Primary Electromagnetic Coupling to Spin ½ Infinite Superpositions

Equation (3.5. 16) tells us what is required of spin ½ superpositions to correctly behave as Dirac fermions, allowing us to solve for $\alpha_{EMP}{}^{-1}$ as a function of the coupling ratio χ.

Starting with Eq. (3.5. 16) $\langle n\rangle = \sum_{n=5,6,7.} (c_n c_n{}^*)n \approx \sqrt{36(1.0045062...)} \approx 6.01350345$.

$$5c_5{}^*c_5 + 6c_6{}^*c_6 + 7c_7{}^*c_7 \approx 6.01350345723 \quad \text{but} \quad 6c_5{}^*c_5 + 6c_6{}^*c_6 + 6c_7{}^*c_7 = 6$$

$$\text{Thus } c_7{}^*c_7 - c_5{}^*c_5 \approx 0.01350340345273$$

As $c_7{}^*c_7 + c_5{}^*c_5 = 1 - c_6{}^*c_6$ we can solve for $c_7{}^*c_7$ & $c_5{}^*c_5$ in terms of $c_6{}^*c_6$.

$$c_7{}^*c_7 \approx 0.50675172 - \frac{c_6{}^*c_6}{2} \quad \& \quad c_5{}^*c_5 \approx 0.49324827 - \frac{c_6{}^*c_6}{2} \qquad (4.1.\ 1)$$

From Eq. (2.3. 12) the Q^2A^2 required to produce this superposition with amplitudes c_n is

$$Q^2A^2 = \sum_{n=5,6,7} c_n * c_n \frac{n^4\hbar^2k^4r^2}{81} \text{ and using Eq.(4.1. 1)}$$

$$\sum_{n=5,6,7} c_n * c_n \; n^4 = 625c_5 * c_5 + 1296c_6 * c_6 + 2401c_7 * c_7 \approx 1524.991 - 217c_6 * c_6$$

Thus $Q^2A^2 = \sum_{n=5,6,7} c_n * c_n \frac{n^4\hbar^2k^4r^2}{81} \approx [18.82705 - 2.67901c_6 * c_6]\hbar^2k^4r^2$ is *the required vector potential squared to produce this spin ½ superposition*. From Eq. (2.2. 4) with $s =$ ½ & $N = 1$ for massive fermions $Q^2A^2 = \frac{2\left[8+8\sqrt{\alpha_{EMP}}\right)\Big]^2}{3\pi}\hbar^2k^4r^2$ *is the available* Q^2A^2.

Equating *required and available*: $2\left[8+8\sqrt{\alpha_{EMP}}\right)\Big]^2 \approx 3\pi[18.82705 - 2.67901c_6 * c_6]$

$$\left[1+\sqrt{\alpha_{EMP}}\right)\Big]^2 \approx [1.386256 - 0.197258c_6 * c_6]$$

$$\alpha_{EMP} \approx \left[\sqrt{1.386256 - 0.197258c_6 * c_6} - 1\right]^2 \qquad (4.1.\ 2)$$

From Eq's. (3.3. 23) & (3.3. 24) $c_6 * c_6(1 - c_6 * c_6) = \sqrt{2/\chi} = 6\sqrt{2/\chi_{EM}}$, we can now solve for α_{EMP} as a function of either χ_{EM} or χ. We then use Eq. (3.3. 24) again to get $\alpha_{EMS}{}^{-1} @ \approx L_P$.

Now both χ_{EM} and χ are fundamentally the same ratio differing only by 36:1, because electron superpositions have six primary charges whereas we define them as one fundamental charge (section 3.3.1) and quarks have only one colour charge (Table 2.2. 1). Because $\chi = \alpha_3^{-1}$ at the cutoff near L_P it is more convenient to work with.

From Eq. (3.3. 23) $c_6 * c_6 = \frac{1}{2} \pm \frac{1}{2}\sqrt{1 - 4\sqrt{\frac{2}{\chi}}}$ and there are two solutions for each χ; one with $c_6 * c_6$ dominant and two smaller $c_5 * c_5$ & $c_7 * c_7$ side modes, or the reverse with $c_6 * c_6$ the minor player and two larger $c_5 * c_5$ & $c_7 * c_7$ side modes. As the values for α_{EMP} with $c_6 * c_6$ dominant fit the Standard Model best, we include only these. Table 4.1. 1 shows these results for $\chi = \alpha_3^{-1}$ at possible cutoffs, but in the range $\chi = 50 \to 51$ only, as we will see in section 4.3 and Figure 4.1. 2 that this range fits reasonably well with the Standard Model. The yellow highlighted row is approximately in the centre of the intersecting region of Figure 4.1. 2.

Coupling Ratio χ	$c_6 * c_6$	$\alpha^{-1}{}_{EM\,Primary}$	$\alpha^{-1}{}_{EMSecondary}$ @ Cutoff
50.00	≈ 0.723607	≈ 75.4414	≈ 104.7798
50.20	≈ 0.724497	≈ 75.5447	≈ 105.3429
50.40	≈ 0.725378	≈ 75.6472	≈ 105.9060
50.60	≈ 0.726250	≈ 75.7488	≈ 106.4692
50.80	≈ 0.727115	≈ 75.8497	≈ 107.0324
51.00	≈ 0.727970	≈ 75.9499	≈ 107.5956

Table 4.1. 1 Coupling ratio χ versus $\alpha^{-1}{}_{EMSecondary}$ in the range $\chi = 50 \rightarrow 51$.

4.1.1 Comparing this with the Standard Model

In the real world of secondary interactions the Standard Model splits the electromagnetic force into the two components $\alpha_1 \& \alpha_2$ for energies greater than the mass/energy of the Z_0 boson or $\approx 91.2\, GeV$. [12]. We want to however compare the Standard Model electromagnetic coupling with the values derived above at the cutoff near the Planck length.

The weak force split obeys $\alpha_{EM}{}^{-1} = \frac{5}{3}\alpha_1{}^{-1} + \alpha_2{}^{-1}$ (4.1. 3)

Also $\alpha^{-1}{}_1 = \frac{3}{5}\alpha_{EM}{}^{-1}Cos^2\theta_W$ & $\alpha_2{}^{-1} = \alpha_{EM}{}^{-1}Sin^2\theta_W$ where θ_W is the Weinberg angle.

Assuming three families of fermions and one Higgs field the SM [13]. predicts

$$\begin{aligned}
\alpha_1{}^{-1} &\approx 58.98 \pm 0.08 - \frac{4.1}{2\pi}\log_e\frac{Q}{91.2} \\
\alpha_2{}^{-1} &\approx 29.6 \pm 0.04 + \frac{3.16666}{2\pi}\log_e\frac{Q}{91.2} \\
\alpha_3{}^{-1} &\approx 8.47 \pm 0.22 + \frac{7}{2\pi}\log_e\frac{Q}{91.2}
\end{aligned} \qquad (4.1.\ 4)$$

Combining Eq's. (4.1. 3) & (4.1. 4)

$$\alpha_{EM}{}^{-1} \approx 127.9 \pm 0.173 - \frac{3.66666}{2\pi}\log_e\frac{Q}{91.2} \qquad (4.1.\ 5)$$

Figure 4.1. 1 plots these four inverse coupling constants. Figure 4.1. 2 plots the intersection of $\alpha^{-1}{}_{EMSecondary}$ predicted in Table 4.1. 1 and the Standard Model prediction for $\alpha^{-1}{}_{EM}$ in Eq. (4.1. 5). It would initially seem in Figure 4.1. 2 that there is an unusually large error band in the predicted results. However $\Delta\alpha^{-1}{}_{EMSecondary} / \Delta\chi^{-1} \approx 2.8$ is approximately constant in this table and the error band in the Standard Model colour coupling $\alpha_3{}^{-1}$ of ± 0.22 in Eq's (4.1. 4)

translates into the much larger error band for $\alpha^{-1}{}_{EMSecondary}$, of $\pm 0.22\times 2.8\approx \pm 0.62$ in Figure 4.1. 2. In section 4.3 we look at superposition cutoffs exponentially tailing off into the Planck region after we have introduced the effect of gravity into our equations.

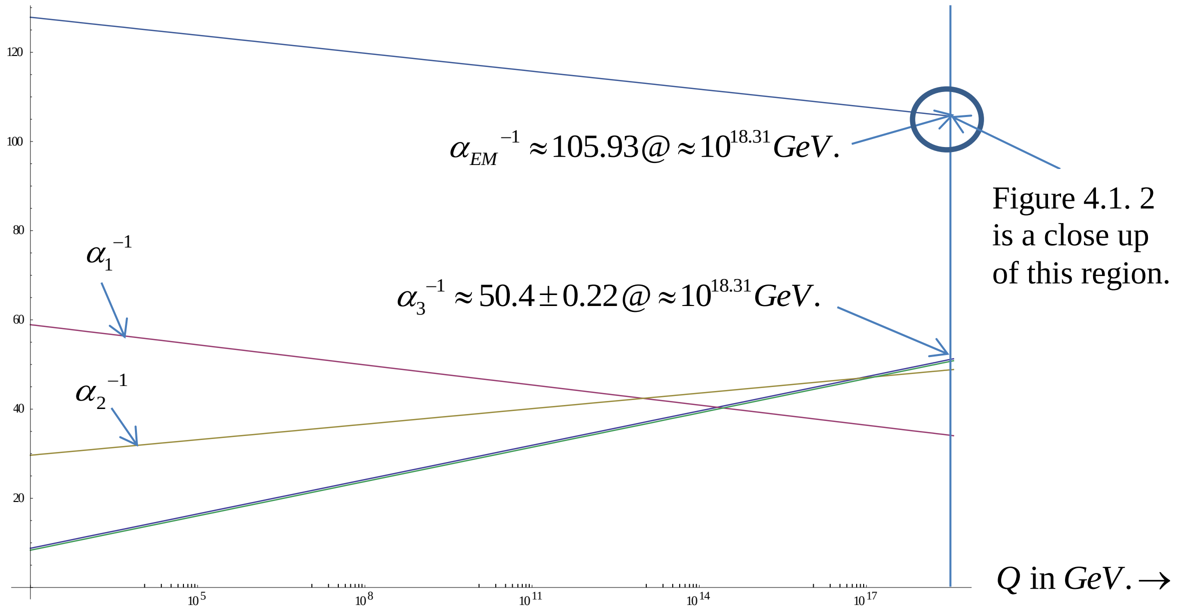


Figure 4.1. 1 Standard Model based on three families of fermions and one Higgs field.

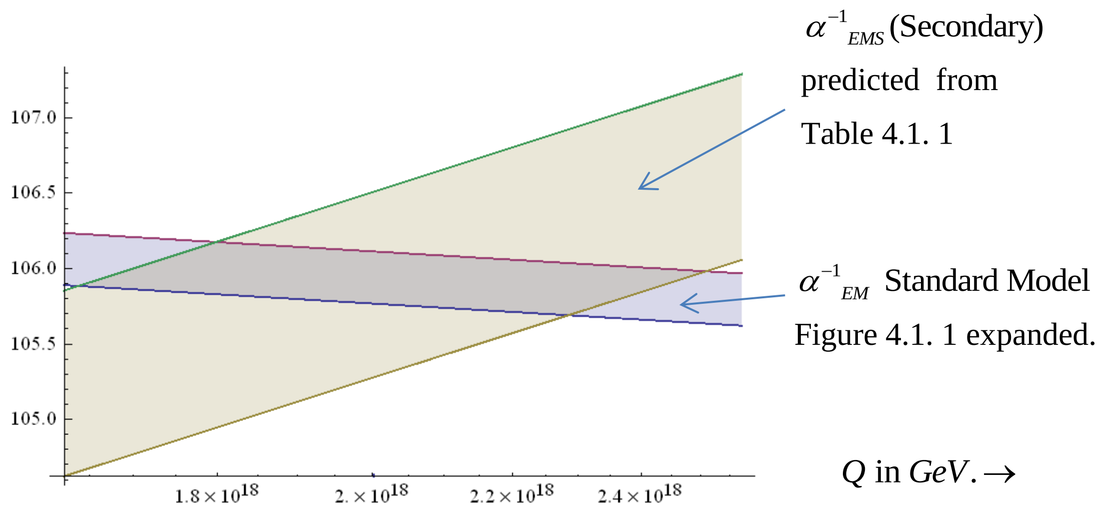


Figure 4.1. 2 Predicted range for $\alpha^{-1}{}_{EM}$ (Bare) $\approx 105.93\pm 0.25$ in the intersecting (or cutoff) region of the Standard Model and this paper.

4.2 Introducing Gravity into our Equations

4.2.1 Initially looking at square superposition cutoffs only

In section 3.2 it was easier to work with one wavefunction ψ_{nk} at a time, and we found later that a superposition of them did not affect our final results; we will do the same here. We also found in Eq's. (3.2. 3) & (3.2. 6) that the integrals for both angular momentum and rest masses are of identical form. They both ended up including the term

$\left[\frac{-1}{1+K_{nk}{}^2}\right]_0^\infty$ which if $K_{nk}cutoff < \infty$ becomes $\left[\frac{-1}{1+K_{nk}{}^2}\right]_0^{K_{nk}cutoff}$ which is equal to

$$1-\frac{1}{1+K_{nk}{}^2 cutoff}=\frac{K_{nk}{}^2 cutoff}{1+K_{nk}{}^2 cutoff}=\frac{1}{1+1/K_{nk}{}^2(cutoff)}=\frac{1}{1+\varepsilon} \tag{4.2. 1}$$

where using Eq. (3.1. 11) the infinitesimal $\varepsilon=\frac{1}{K_{nk}{}^2 cutoff}=\frac{2m_0{}^2c^2}{n^2\hbar^2(k_{cutoff})^2 s}$ (4.2. 2)

For integral or half integral $\hbar$ angular momentum absolute precision is required but Eq. (3.2. 6) now gives us $\mathbf{L}_z(Total)=\frac{sm\hbar}{2}\left[\frac{-1}{1+K_{nk}{}^2}\right]_0^{K_{nk}cutoff}=\frac{sm\hbar}{2}\frac{1}{1+\varepsilon}$. So can the effect of gravity increase our probabilities from $sN\cdot\frac{dk}{k}$ to $sN\cdot(1+\varepsilon)\frac{dk}{k}$? We will initially address only *massive infinite superpositions* where $N=1$ in Eq.(2.2. 4).

The first question we need to address is what is the effective preon mass to be used when coupling to gravity? In Eq. (3.1. 4) we said the preon rest mass is $m_0/(8\gamma_{nk}\sqrt{2s})$ for each of the 8 preons that build a spin ½ particle of rest mass m_0. Now gravity couples to the total mass including the kinetic energy. It also couples to pressure in Einstein's energy-momentum tensor, but to keep things simple initially, we will temporarily ignore possible pressure terms. At the start of the interaction each preon mass is $m_0/(8\gamma_{nk}\sqrt{2s})$ and after the interaction (Figure 3.1. 3) it is $m_0\gamma_{nk}(1+\beta_{nk}{}^2)/(8\sqrt{2s})$. Let us think semi classically again and see where it leads us. We have been using magnitudes of velocities as they are the most convenient way to express our equations even if not the conventional language of quantum mechanics. The interaction with the zero point fields takes the momentum of each preon from zero to $2m_0\gamma_{nk}\beta_{nk}c/(8\sqrt{2s})$ (Figure 3.1. 3). While this happens as a quantum step change let us imagine it as a virtually infinite acceleration from zero velocity to $2\beta_{nk}/(1+\beta_{nk}{}^{2})$, which is the relativistic velocity addition (see Figure 3.1. 1) of 2 equal steps of β_{nk}. At the half way point after one step the velocity is β_{nk} (the velocity of the CMF, the preon mass has increased to $m_0/(8\sqrt{2s})$. We can imagine this as being like the central point of a quantum interaction.

We will conjecture this midway point preon mass $m_0/(8\sqrt{2s})$ is the mass value that gravity acts on and we will see that it is indeed the only value that makes all the equations fit. Also it does not make sense to choose either of the end point masses. We can also get reassurance from the properties of the Feynman transition amplitude which tells us in Eq. (3.1. 15)

$\frac{(p_i+p_f)^z}{(p_i+p_f)^0}=\frac{2m_0\gamma_{nk}\beta_{nk}}{2m_0\gamma_{nk}}=\beta_{nk}$ and the ratio of space to time polarization in the LF is $\beta_{nk}{}^{2}$.

This centre of momentum velocity tells us the key properties of the interaction. (Still ignoring coupling to this momentum term.) We will thus assume we have 8 preons in each ψ_{nk} of effective gravitational mass $m_0/(8\sqrt{2s})$ with effective total gravitational mass $m_0/\sqrt{2s}$. To put the gravitational constant in the same form as the other coupling constants we need to divide it by $\hbar c$. The gravitational coupling amplitude is thus $m_0\sqrt{G_P/(2s\hbar c)}$ to the gravitational zero point field, where G_P is the primary gravitational constant. Now this gravitational amplitude can be regarded as a complex vector just as colour and electromagnetism. We assumed for simplicity, as they are both spin 1 field particles, that colour and electromagnetism were parallel. Gravity is spin 2 so it can be expected that it may well be at a different complex angle to the other two. In fact if we put it at right angles to the other two the equations all seem to fit. We will thus conjecture using Eq.(3.3. 24) that (temporarily ignoring possible gravitational coupling to any possible pressure component):

The gravitational coupling amplitude is $im_0\sqrt{G_P/(2s\hbar c)}=im_0\sqrt{\chi\cdot G_S/(2s\hbar c)}$ (4.2. 3)

Because the actual cutoff (see Figure 4.3. 1) is well inside the Planck length we will *put the secondary gravitational coupling constant G_s to a bare Planck mass equal to the measured gravitational constant G.* So modifying Eq's. (2.2. 1) to (2.2. 3) by adding Eq.(4.2. 3)

$$Q^2A^2=\left[\frac{\left|8+8\sqrt{\alpha_{EMP}}+im_0\sqrt{\chi\cdot G/(2s\hbar c)}\right|^2}{3\pi sN}\hbar^2k^4r^2\right]\cdot\left[\frac{sN\cdot dk}{k}\right]$$

$$Q^2A^2 = \left[\frac{\left[8+8\sqrt{\alpha_{EMP}})\right]^2}{3\pi sN} \hbar^2 k^4 r^2 \right] . \left[\frac{sN(1+\varepsilon')dk}{k} \right] \text{ where } \varepsilon' = \frac{m_0{}^2 \chi \cdot G}{2s\hbar c(8+8\sqrt{\alpha_{EMP}})^2}$$

Our previous wavefunctions ψ_k required $Q^2A^2 = \dfrac{\left[8+8\sqrt{\alpha_{EMP}})\right]^2}{3\pi sN} \hbar^2 k^4 r^2$ from Eq.(2.2. 4).

Gravity can increase the probability of our previous wavefunctions ψ_k by $1+\varepsilon'$ as required to obtain precision in our integrals for $\hbar/2\,\&\,\hbar$ if $K_{nk}cutoff < \infty$.

Using Eq.(4.2. 2) now put $\varepsilon' = \dfrac{m_0{}^2 \chi \cdot G}{2s\hbar c(8+8\sqrt{\alpha_{EMP}})^2} = \varepsilon = \dfrac{1}{K_{nk}{}^2 cutoff} = \dfrac{2m_0{}^2 c^2}{sn^2\hbar^2 (k_{cutoff})^2}$ (4.2. 4)

4.2.2 Including a Pressure term from Einstein's Energy Momentum Tensor

To keep the virtual particles in quantum mechanical orbits there is a pressure term $\propto \beta_{nk}{}^2$ and the gravitational coupling amplitude Eq. (4.2. 3) becomes $im_0(1+\beta_{nk}{}^2)\sqrt{\chi \cdot G/(2s\hbar c)}$, and ε' becomes $\varepsilon'' \to \varepsilon'(1+\beta_{nk}{}^2)^2$ where β_{nk} is the velocity at the central point of this interaction which varies with *k*. To find an averaged effect for this pressure term we need to integrate a varying ε''. Using the equations leading up to Eq. (3.2. 6) these integrals become:

$$\mathbf{L}_z(Total) = sm\hbar \int_0^{K_c} \frac{K_{nk}{}^2}{(1+K_{nk}{}^2)^2} \frac{dK_{nk}}{K_{nk}} [1+\varepsilon''] = sm\hbar \int_0^{K_c} \frac{K_{nk}{}^2}{(1+K_{nk}{}^2)^2} \frac{dK_{nk}}{K_{nk}} [1+\varepsilon'(1+\beta_{nk}{}^2)^2]$$

From Eq. (3.1. 12) $\beta_{nk}{}^2 = \dfrac{K_{nk}{}^2}{1+K_{nk}{}^2}$ and with appropriate factors the averaged effect on ε'' is:

$$\varepsilon'' \to \varepsilon' \int_0^{K_c} \frac{2K_{nk}{}^2}{(1+K_{nk}{}^2)^2} \frac{dK_{nk}}{K_{nk}} [1+\frac{K_{nk}{}^2}{(1+K_{nk}{}^2)}]^2 \text{ or } \varepsilon'' \to \frac{7}{3}\varepsilon' \text{ as } K_{cutoff} \to \infty \qquad (4.2.\ 5)$$

Equation (4.2. 4) becomes $\varepsilon'' \approx \dfrac{(7/3)m_0{}^2 \chi \cdot G}{2s\hbar c(8+8\sqrt{\alpha_{EMP}})^2} \approx \varepsilon = \dfrac{1}{K_{nk}{}^2 cutoff} = \dfrac{2m_0{}^2 c^2}{sn^2\hbar^2 (k_{cutoff})^2}$

or $\dfrac{(7/3)\chi \cdot G}{4\hbar c(8+8\sqrt{\alpha_{EMP}})^2} \approx \dfrac{c^2}{n^2\hbar^2 (k_{cutoff})^2}$ and $\dfrac{(7/3)\chi}{256(1+\sqrt{\alpha_{EMP}})^2} \dfrac{G\hbar}{c^3} \approx \dfrac{1}{n^2 (k_{cutoff})^2}$

$$\text{But } L_P{}^2 = \frac{G\hbar}{c^3} \quad \text{and} \quad \frac{(7/3)\chi \cdot L_P{}^2}{256(1+\sqrt{\alpha_{EMP}})^2} \approx \frac{1}{n^2(k_{cutoff})^2} \tag{4.2. 6}$$

Now consider the radial probability of $\psi_{nk} = C_{nk} r^3 \exp(-n^2k^2r^2/18)Y(\theta,\varphi)$.

Ignoring factors C_{nk} & $Y(\theta,\varphi)$ and differentiating $P_r = r^8 \exp(-n^2k^2r^2/9)$ we get

$$\frac{dP_r}{dr} = \left[8 - \frac{2n^2k^2r^2}{9}\right].\left[r^7 \exp(-n^2k^2r^2/9)\right] = 0 \text{ when } r^2 = \frac{36}{n^2k^2} \text{ or } r_{peak} = \frac{6}{nk}$$

$$\text{Thus for a superposition} \quad \langle r_{peak} \rangle = \frac{6}{\langle n \rangle k} \quad \& \quad \& \quad \langle r_{peak} \rangle^2 = \frac{36}{\langle n \rangle^2 k^2} \tag{4.2. 7}$$

Also for superpositions using Eq. (3.2. 10) , Eq. (4.2. 2**)** becomes:

$$\varepsilon = \frac{1}{\langle K_{nk} cutoff \rangle^2} = \frac{2m_0{}^2c^2}{s\langle n \rangle^2 \hbar^2 (k_{cutoff})^2} \tag{4.2. 8}$$

Putting equations (4.2. 6), (4.2. 7) & (4.2. 8) together

$$\frac{(7/3)\chi \cdot L_P{}^2}{256(1+\sqrt{\alpha_{EMP}})^2} \approx \frac{1}{\langle n \rangle^2 (k_{cutoff})^2} = \frac{\langle r_{peak@cutoff} \rangle^2}{36}. \tag{4.2. 9}$$

The expectation value of the maximum probability radius at this cutoff is independent of superposition value n. In section 4.1, from Eq's. (4.1. 2) we find $1+\sqrt{\alpha_{EMP}} \approx 1.115$ and thus:

$$\frac{\langle r_{peak@cutoff} \rangle}{L_P} \approx \frac{\sqrt{\chi}}{1.946} \tag{4.2. 10}$$

The peak radial probability at cutoff depends only on ratio χ for all massive $N=1$ fundamental particles represented as infinite superpositions. From Figure 4.1. 1 & Figure 4.1. 2 we find $\chi = \chi_C = \alpha_3^{-1} \approx 50.4$ at the cutoff and putting this into Eq. (4.2. 10) we get

$$r_{peak@cutoff} = 1/k_{cutoff} \approx 3.65 L_P. \tag{4.2. 11}$$

We have no idea what might happen inside the Planck region, so cutoffs outside the Planck length make intuitive sense. Figure 4.2. 1 illustrates the radial Gaussian probability distribution of such a square cutoff wavefunction just outside the Planck length. The Planck energy $E_P \approx 1.22\times10^{19}GeV$. and using Equ's. (4.2. 9) & (4.2. 11):

$$\text{Assuming a square cutoff:} \quad \hbar k_{cutoff}c \approx 3.34\times10^{18}GeV. \tag{4.2. 12}$$

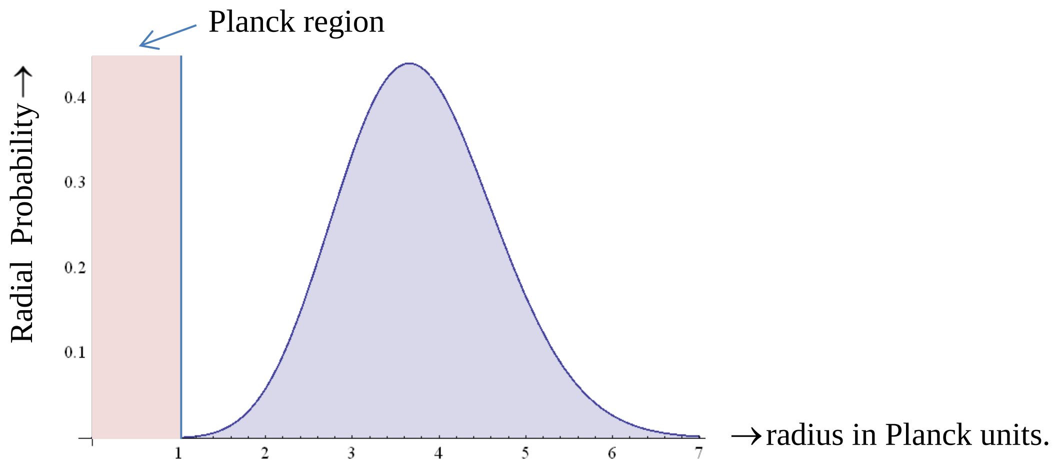


Figure 4.2. 1 Radial probability for square cutoff wavefunction ψ_k @$\approx 3.34\times10^{18}GeV$.

In the above we have only analysed massive $N=1$ superpositions. To analyse $N=2$ superpositions we need the results of section 6.2.

4.3 Exponential Superposition Cutoffs

Because a square cutoff is close to the Planck length (Figure 4.2. 1) it is reasonable to assume that any cutoff should be as sharp as possible to minimize penetration of the Planck region. Exponentials of the type $(-k^2cutoff\,/\,k^2)$ cut off faster than exponential $(-k(cutoff)/k)$. Our wavefunctions for superpositions are also squared exponentials, so we will try this possibility.

Using $\int_0^\infty Exp\left(-\frac{\langle K_{nk(cutoff)}\rangle^2}{\langle K_{nk}\rangle^2}\right)\frac{2d\langle K_{nk}\rangle}{\langle K_{nk}\rangle^3} = \frac{1}{\langle K_{nk(cutoff)}\rangle^2}$ & $\langle K_{nk(cutoff)}\rangle >>>> 1$ we can show:

$$\int_0^{K_{nk}(Cutoff)} \frac{2\langle K_{nk}\rangle}{(1+\langle K_{nk}\rangle^2)^2} d\langle K_{nk}\rangle \approx \int_0^\infty \left[1 - Exp\frac{-\langle K_{nkCutoff}\rangle^2}{\langle K_{nk}\rangle^2}\right]\frac{2\langle K_{nk}\rangle}{(1+\langle K_{nk}\rangle^2)^2} d\langle K_{nk}\rangle \tag{4.3. 1}$$

Both produce the result : $\approx 1-\dfrac{1}{K^2_{\ nkCutoff}}$ but only if $K_{nkCutoff} \to \infty$. Also from Eq's (3.2. 10)

$$\frac{\left\langle K_{nkCutoff}\right\rangle}{\left\langle K_{nk}\right\rangle}=\frac{k_{Cutoff}}{k} \qquad \text{and} \qquad 1-Exp\frac{-\left\langle K_{nkCutoff}\right\rangle^2}{\left\langle K_{nk}\right\rangle^2} \quad = \quad 1-Exp\frac{-k_{Cutoff}{}^2}{k^2} \tag{4.3. 2}$$

Where $\left\langle K_{nkCutoff}\right\rangle \& \, k_{Cutoff}$ are the same as for the square cutoff in Eq. (4.2. 8). Thus an *exponential cutoff of this form gives identical results to the original square cutoff for both angular momentum and rest masses* as in section 4.2.1.

Modifying Eq. (3.3. 15) for the *interaction probability* of massive $N=1$ spin 1 bosons between two massive $N=1$ spin ½ superpositions

$$\frac{\left[2s_{1/2}N_1c_{6a}*c_{6a}(1-c_{6a}*c_{6a})\right]^2\left[2s_1N_1c_{5b}*c_{5b}(1-c_{5b}*c_{5b})\right]^2}{q^4} \text{ for } N=1 \text{ bosons} \tag{4.3. 3}$$

We used superposition probabilities $s_{1/2}N_1dk/k,\ s_1N_2dk/k$ from Eq. (2.2. 4) in deriving Eq. (3.3. 15) but in Eq. (4.3. 3) we used $s_{1/2}N_1dk/k \,\&\, s_1N_1dk/k$ as $N=1$ in all superpositions here. (These derivations in section 3.3 were based on a thought experiment at very large spacings to keep things simple but the ratios must still apply at all wavenumbers.) So if all superposititions cutoff as in Eq. (4.3. 2) the *interaction probability* between all pairs of spin ½ superpositions exchanging massive $N=1$ spin 1 superposition bosons should become near the Planck length

$$\frac{\left[2s_{1/2}N_1c_{6a}*c_{6a}(1-c_{6a}*c_{6a})\right]^2\left[2s_1N_1c_{5b}*c_{5b}(1-c_{5b}*c_{5b})\right]^2}{q^4}\left[1-Exp\frac{-k^2Cutoff}{k^2}\right]^4 \tag{4.3. 4}$$

Thus we can conjecture with this type of exponential cutoff that near the Planck length

$$\text{Interaction probabilities between pairs of superpositions is} \propto \left[1-Exp\frac{-k^2Cutoff}{k^2}\right]^4 \tag{4.3. 5}$$

Where $k(cutoff)=k(square\ cutoff)\approx 3.34\times 10^{18}GeV$ as in Eq.(4.2. 12)

Figure 4.3. 1 plots these possible exponential cutoffs. It includes the superimposed high energy interaction cutoff region from Figure 4.1. 2 which is at a slightly lower energy but very close. Of course we have only conjectured a possible exponential cutoff here however different types of exponential cutoffs are all very close to this region. The cutoff region in

Figure 4.1. 2 ranges from $\approx 1.6\times10^{18} GeV$ to $\approx 2.5\times10^{18} GeV$ centered at $\approx 2.04\times10^{18} GeV$. To solve superposition amplitudes we will use this central cutoff energy for the coupling ratio

$$\chi = \alpha_3^{-1} \approx 50.4 \pm 0.4 @ \approx 10^{18.31} GeV. \tag{4.3. 6}$$

Interaction cutoff region suggested by Figure 4.1. 2

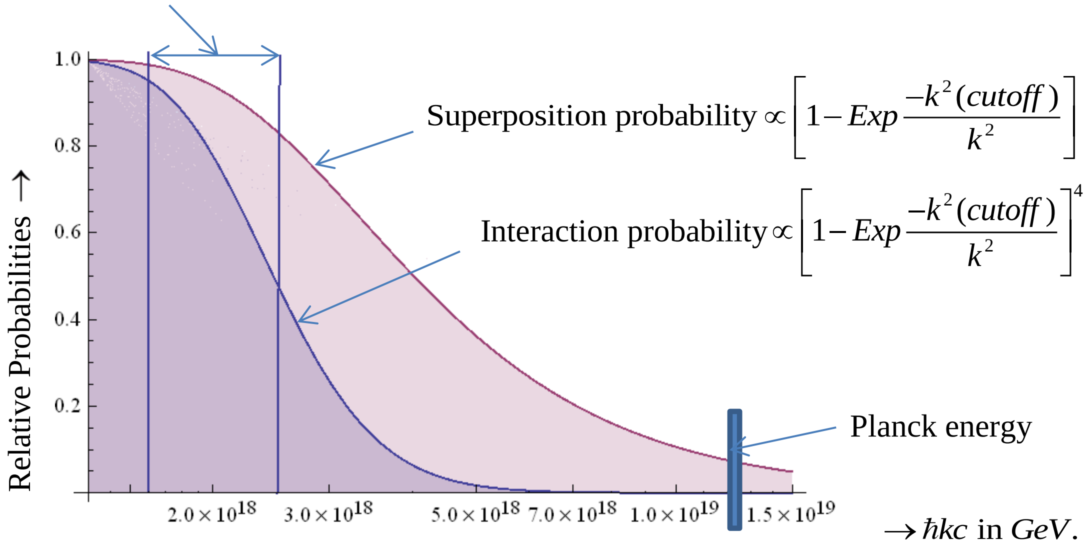


Figure 4.3. 1 Exponential cutoffs where $\hbar k(cutoff)c \approx 3.34\times10^{18} GeV$. see Eq. (4.2. 12).

4.4 Solving for spin ½, spin 1 and spin 2 superpositions

Superpositions with $N = 2$ are covered in section 6.2 but Eq. (4.3. 6) and Eq. (3.3. 23) extended by keeping $N \cdot s$ constant as in Eq. (4.4. 1) allow us to solve various combinations of spins ½, 1 or 2 and $N = 1$ or $N = 2$.

$$\begin{array}{cccc} (N=2)\times(\text{Spin }2) & (N=2)\times(\text{Spin }1) & (N=1)\times(\text{Spin }1) & (N=1)\times(\text{Spin }\tfrac{1}{2}) \\ & \text{or } (N=1)\times(\text{Spin }2) & \text{or } (N=2)\times(\text{Spin }\tfrac{1}{2}) & \end{array}$$

$$8c_{4c}*c_{4c}(1-c_{4c}*c_{4c}) = 4c_{4b}*c_{4b}(1-c_{4b}*c_{4b}) = 2c_{5b}*c_{5b}(1-c_{5b}*c_{5b}) = c_{6a}*c_{6a}(1-c_{6a}*c_{6a})$$
$$= \sqrt{2/\chi} \approx \sqrt{2/50.4} \approx 0.1992 \tag{4.4. 1}$$

Starting with spin ½ we can solve this to get $c_6 * c_6 \approx 0.7254$ as the dominant value. Putting $c_6 * c_6 \approx 0.7254$ into Eq. (4.1. 2) or alternatively using Table 4.1. 1

$$\alpha_{EMP} \approx \left[\sqrt{1.386256 - 0.197258 c_6 * c_6} - 1\right]^2 \approx 75.65^{-1} \tag{4.4. 2}$$

From Eq. (2.2. 4) the available $Q^2A^2 = \frac{\left[8+8\sqrt{\alpha_{EMP}})\right]^2}{3\pi sN}\hbar^2k^4r^2$ with probability $\frac{sN \cdot dk}{k}$ where we ignore the infinitesimal factor of $(1+\varepsilon)$ due to gravity. And from Eq. (2.3. 12)

$$Q^2A^2 = \sum_n c_n * c_n \frac{n^4\hbar^2k^4r^2}{81} = \frac{\left[8+8\sqrt{\alpha_{EMP}})\right]^2}{3\pi sN}\hbar^2k^4r^2$$

$$\begin{aligned} \sum c_n * c_n \cdot n^4 &\approx 1367.58 \text{ for (spin } 1/2 \times N = 1) \\ &\approx 683.79 \text{ for (spin } 1 \times N = 1) \text{ or (spin } 1/2 \times N = 2) \\ &\approx 341.9 \text{ for (spin } 1 \times N = 2) \text{ or (spin } 2 \times N = 1) \\ &\approx 170.95 \text{ for (spin } 2,\ N = 2). \end{aligned} \tag{4.4. 3}$$

The same primary electromagnetic coupling α_{EMP} builds all fundamental particles, allowing Eq. (4.4. 3) to be true. Using Eq's. (4.4. 1), (4.4. 3) & $\sum_n c_n * c_n = 1$ we get Table 4.3. 1.

Mass Type	Spin	N	$c_3 * c_3$	$c_4 * c_4$	$c_5 * c_5$	$c_6 * c_6$	$c_7 * c_7$
Infinitesimal mass gravitons	2	2	0.8173	0.0256	0.1571		
Infinitesimal & Massive bosons	1	2	0.4847	0.0526	0.4627		
Massive (dark matter?) gravitons	2	1	0.4847	0.0526	0.4627		
Massive (real?) bosons	1	1		0.0134	0.8878	0.0988	
Massive (virtual?) fermions	½	2		0.0134	0.8878	0.0988	
Massive (real?) fermions	½	1			0.1305	0.7254	0.1441

Table 4.3. 1 Approximate probabilities for various possible superpositions.

Table 4.3. 1 is approximate as it uses the midpoint of the energy range in Figure 4.1. 2 also three member superpositions seem to fit with the Standard Model. A massive $N = 1$ spin 2 graviton type Dark Matter possibility interacting only with $N = 2$ spin 2 gravitons is included. There are a wide range of possibilities in this table (hence the question marks) that can only be more fully covered when we get to section 6.2 and the following subsections.

To this point this paper has tried to demonstrate that infinite superpositions can behave as their Standard Model equivalents. There may well be many errors in the details of derivations; but if the principles behind it all are at least on the right track and fundamental particles can be built by borrowing energy and mass from the zero point fields, then as we will see in what follows this can have some very significant consequences.

5 The Expanding Universe and General Relativity

5.1 Zero point energy densities are limited

If the fundamental particles can be built from energy borrowed from zero point fields and as this energy source is limited, (particularly at very long wavelengths) there must be implications for the maximum possible densities of these fundamental particles. In section 2.2.3 we discussed how the preons that build the fundamental particles are born from a Higg's type scalar field with zero momentum in the laboratory rest frame. In this frame they have an infinite wavelength and can thus be borrowed from anywhere in the universe. This would suggest that there should be little effect on localized densities, but possibly on overall average densities in any or all of these universes. So which fundamental particle is there likely to be most of? *Working in Planck, or natural units with* $G = 1$, the gravitational coupling constant between Planck masses is one. As an example there are approximately $M \approx 10^{61}$ Planck masses within our causally connected region of the universe. They have an average distance between them of approximately the radius R of this region. Thus there must be approximately $M^2 \approx 10^{122}$ virtual gravitons with wavelengths of the order of radius R within this same volume. No other fundamental particle is likely to approach these values, for example the number of virtual photons *of this extreme wavelength* is much smaller. (Virtual particles emerging from the vacuum are covered in section 6.2.1.) If this density of virtual gravitons needs to borrow more energy from the zero point fields than what is available at these extreme wavelengths does this somehow control the maximum possible density of any causally connected universe? There may well be insufficient long wavelength virtual gravitons to allow gravity to apply at vast cosmic scales such as the distance to the horizon. At smaller scales within these causally connected regions, where the wavelengths of the virtual gravitons involved are shorter, and the zero point energy densities (which vary as $k^3 dk$ or $d\lambdabar / \lambdabar^5$) are much greater, gravity would apply in the usual manner.

5.1.1 Virtual graviton density at wavenumber *k* in a causally connected Universe

From here on we will work in natural or Planck units where $\hbar = c = G = 1$.

General Relativity predicts nonlinear fields near black holes, but in the low average densities of typical universes *we can assume approximate linearity*. We can also ignore momentum as the average local velocities are low (i.e. $\beta << 1$). We should also be able to simply apply the equations in sections 3.3.4 & 3.5 to spin 2 virtual graviton emissions, as they should apply equally to both spins 1 & 2. We will assume spherically symmetric $l = 3$ wavefunctions emit both spin 1 & 2 scalar virtual bosons, and $l = 3, m = \pm 2$ states emit both $m = \pm 1$, & $m = \pm 2$

bosons. Section 3.3.4 derived the electrostatic energy between infinite superpositions. In flat space we looked at the amplitude that each equivalent point charge emits a virtual photon, and then focused on the interaction terms between them. Thus we can use the same scalar wavefunctions Eq's. (3.4. 1) for virtual scalar gravitons as we did for virtual photons. For virtual photons in section 3.4.1 we showed that using

$$(\psi_1+\psi_2)^*(\psi_1+\psi_2)=\psi_1{}^*\psi_1+\psi_1{}^*\psi_2+\psi_2{}^*\psi_1+\psi_2{}^*\psi_2$$

the interaction term is $\psi_1{}^*\psi_2+\psi_2{}^*\psi_1=\dfrac{4k}{4\pi r_1 r_2}e^{-k(r_1+r_2)}\cos k(r_1-r_2)$ (5.1. 1)

This equation is strictly true only in flat space but it is still approximately true if the curvature is small or when $2m/r \lll\lll 1$*, which we will assume applies almost everywhere throughout the universe except in the infinitesimal fraction of space close to black holes*. In both sections 3.3.4 & 3.5, for simplicity and clarity, we delayed using coupling constants and emission probabilities in the wavefunctions until necessary. We will do the same here. There will also be some minimum wavenumber *k* which *we will call* $k_{\min}$ where for all $k<k_{\min}$ there will be insufficient zero point energy (or mode) density available. We need this equation to apply for all values of *k* down to this minimum value which we find is $k_{\min}\approx 1/R_{CausallyConnectedHorizon}$. At these extreme values the infinitesimal $N=2$ boson rest mass must be included in the exponential term, as it is also $m'\approx 1/R_{CCH}$. (This effect is only significant at very large radii.) In section 6.2 we will look more closely at $N=2$ infinitesimal rest masses finding that $\left\langle K_{nk\min}{}^2\right\rangle\approx 1$ and using Eq.(3.1. 11) $\dfrac{s\left\langle n^2\right\rangle k_{\min}{}^2}{2m_0{}^2}\approx 1$.

For spin 2 gravitons $\left\langle K_{nk\min}{}^2\right\rangle=\dfrac{\left\langle n^2\right\rangle k_{\min}{}^2}{m_0{}^2}\approx 1$ or $m_0{}^2\approx\left\langle n^2\right\rangle k_{\min}{}^2$ (5.1. 2)

The *k* in the exponential term of Eq. (5.1. 1) becomes when $k=k_{\min}$

$$k_{\min}\to\sqrt{k_{\min}{}^2+m_0{}^2}=\sqrt{k_{\min}{}^2+\left\langle n^2\right\rangle k_{\min}{}^2}=k_{\min}\sqrt{1+\left\langle n^2\right\rangle} \quad (5.1.\ 3)$$

From Table 4.3. 1 we find for $N=2$ spin 2 gravitons that $\left\langle n^2\right\rangle\approx 11.69$ so that

$$k'_{\min}\to\ \approx k_{\min}\sqrt{1+11.69}\approx 3.56k_{\min} \quad (5.1.\ 4)$$

Thus the interaction term in Eq. (5.1. 1) becomes when $k=k_{\min}$

$$\psi_1 {*} \psi_2 + \psi_2 {*} \psi_1 = \frac{4k_{\min}}{4\pi r_1 r_2} e^{-3.56k_{\min}(r_1+r_2)} \cos k_{\min}(r_1 - r_2) \tag{5.1. 5}$$

In Eq. (5.1. 5) constant $r_1 + r_2$ describe ellipses, and constant $r_1 - r_2$ describe hyperbolae. Using elliptical coordinates in Figure 3.4. 2 to integrate over all space at the constant wavenumber $k_{\min}$, and putting r as the distance between the two Planck masses or charges:

At minimum wavenumber $k_{\min}$; $\iiint (\psi_1 {*} \psi_2 + \psi_2 {*} \psi_1)\ dv = \dfrac{2e^{-3.56k_{\min}r}\sin(k_{\min}r)}{k_{\min}r}$ (5.1. 6)

Now consider *one Planck mass* at any point P somewhere in the interior region of a typical universe, and let the average density be ρ_U (subscript U for homogeneous universe density) Planck masses per unit volume. Consider spherical shells around point P of radius r and thickness dr with $dm = \rho dv = 4\pi r^2 dr \rho_U$. As in section 3.4.1 when we derived Eq. (3.4. 3) we used a scalar emission probability $(2\alpha/\pi)(dk/k)$ which here becomes $(2/\pi)(dk/k)$ between Planck masses where $\alpha = 1$. We want the number of virtual gravitons at P of wavenumber $k_{\min}$ interacting between this point Planck mass and the causally connected universe *where the radius of the causally connected horizon is* R_{CCH}. Using Eq. (5.1. 6) this becomes

$$\left(\frac{2}{\pi}\cdot\frac{dk_{\min}}{k_{\min}}\right)\int_{r=0}^{R_{CCH}} \frac{2e^{-3.56k_{\min}r}\sin(k_{\min}r)}{k_{\min}r}\cdot 4\pi r^2 dr\rho_U = \frac{16\rho_U dk_{\min}}{k_{\min}{}^4}\int_0^{(k_m R_{CCH})} (k_{\min}r)e^{-3.56(k_{\min}r)}\sin(k_{\min}r)d(k_{\min}r) \tag{5.1. 7}$$

Assuming a simple sharp cutoff at the causally connected horizon R_{CCH} this is the virtual graviton coupling between one Planck mass and all the other Planck masses inside the horizon at wavenumber $k_{\min}$. In a homogeneous universe we can carry out this same integral at all points (*at the same cosmic time* T), all at the centre of a causally connected region of the same radius R_{CCH}. To get the total virtual graviton density we just multiply Eq. (5.1. **7**) by $\rho_U / 2$ (so as to not count all pairs of Planck masses twice).

From here on define $\Upsilon = k_{\min} R_{CausallyConnectedHorizon}$ in radians. (5.1. 8)

Thus $\rho_{Gk\min} = \dfrac{8\rho_U{}^2 dk_{\min}}{k_{\min}{}^4}\displaystyle\int_0^{\Upsilon}(k_{\min}r)e^{-3.56(k_{\min}r)}\sin(k_{\min}r)d(k_{\min}r) = \dfrac{8\rho_U{}^2 dk_{\min}}{k_{\min}{}^4}\cdot A$ (5.1. 9)

Where for brevity we have simply put $\displaystyle\int_0^{\Upsilon}(k_{\min}r)e^{-3.56(k_{\min}r)}\sin(k_{\min}r)d(k_{\min}r) = A$

5.2 Relating this to General Relativity

The above assumes that at any cosmic time *T*, there is always some value $k_{\min}$ where the borrowed energy density $E_{Gk} = E_{ZPk}$ the available zero point energy density. *We have also assumed so far that the mass in the universe is like a perfect fluid and homogeneous*, also that space is essentially flat on average. Thus all observers fixed relative to comoving coordinates and at the same cosmic time T must measure the same density of virtual gravitons $\rho_{Gk\min}$ at the minimum value $k_{\min}$, as in Eq**.** (5.1. 9). This equation must be true for all such comoving observers. But what happens if we now change the uniform fluid mass density ρ_U by putting an initially small mass concentration $+m_1$ at some point, because when we are near a mass concentration we would expect it to locally increase this density $\rho_{Gk\min}$? However General Relativity tells us that near mass concentrations the metric changes, radial rulers shrink and local observers measure larger radial lengths. This expands volumes locally and should lower their measurement of the background $\rho_{Gk\min}$. Can these effects balance each other, so that the extra gravitons at $k_{\min}$ generated by a nearby mass concentration can bring this density back to the flat space or background $\rho_{Gk\min}$ of Eq. (5.1. 9)? Changes in the metric will also change the local measurement of $k_{\min} \to k'_{\min} = k_{\min}(1-2m/r)^{-1/2}$ but we will initially only consider the case where $m/r <<<< 1$ and look at the effect of this at the end of section 5.2.2.

5.2.1 Approximations that are necessary with possibly important consequences

Let us refer back to Eq. (3.4. 2) and the steps we took in section 3.4.1 to derive it; but with $k \to k_{\min}$ and $k \to k'_{\min} = 3.56k_{\min}$ in the exponential term of Eq. (3.4. 2)

$$\psi_1 * \psi_2 + \psi_2 * \psi_1 = \frac{4k_{\min}}{4\pi r_1 r_2} e^{-3.56k_{\min}(r_1+r_2)} \cos k_{\min}(r_1 - r_2) \tag{5.2. 1}$$

and assume that *space has to be approximately flat* with errors $\propto 1-(1-2m/r)^{1/2} \approx m/r$. If we now focus on Figure 3.4. 2 , equation (5.2. 1) is the probability that a virtual graviton of wavenumber $k_{\min}$ is at the point *P if all other factors are one*. Let us now put a mass of m_1 Planck masses at the Source 1 point in Figure 3.4. 2 or as in Figure 5.2. 1. Also assume that the *point P is close to mass* m_1 at distance r_1 as in Figure 5.2. 1 and *the vast majority of the rest of the mass inside the horizon* R_{CCH} is at various radii *r* equal to the r_2 of Eq. (5.2. 1) where $r_2 = r >> r_1$. *Only under these conditions can we approximate Eq. (5.2. 1)as*

$$\psi_1 * \psi_2 + \psi_2 * \psi_1 = \frac{4k_{\min}}{4\pi r_1 r} e^{-3.56 k_{\min} r} \cos(-k_{\min} r) \qquad (5.2.\ 2)$$

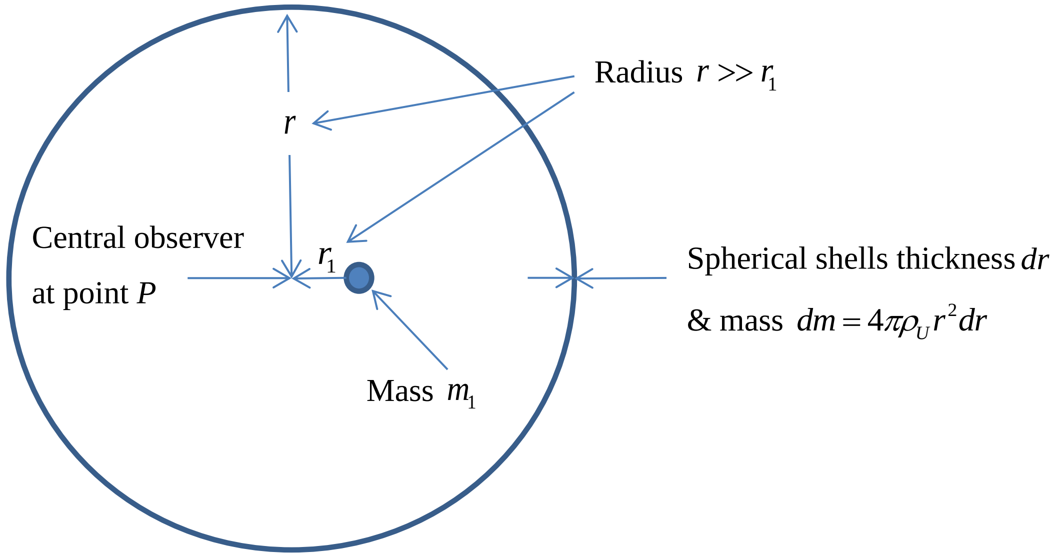


Figure 5.2. 1

As we have assumed average particle velocities are low this is *a scalar interaction* (as in section 3.4.1) and as *there are no directional effects* we can consider simple spherical shells of thickness dr and radius r around a central observer at the point P which have mass $dm = \rho_U 4\pi r^2 dr$. At each radius r the coupling factor $(2\alpha/\pi)(dk/k)$ we used in deriving Eq. (3.4. 3) becomes (as the coupling constant $\alpha = 1$ between Planck masses)

$$\frac{2m_1}{\pi} dm \frac{dk_{\min}}{k_{\min}} = \frac{2m_1}{\pi} \frac{dk_{\min}}{k_{\min}} \rho_U 4\pi r^2 dr \qquad (5.2.\ 3)$$

Including this factor Eq. (5.2. 2) becomes

$$\begin{aligned}(\frac{2m_1}{\pi}\frac{dk_{\min}}{k_{\min}}\rho_U 4\pi r^2 dr)(\psi_1 * \psi_2 + \psi_2 * \psi_1) &\approx \frac{2m_1}{\pi}\frac{dk_{\min}}{k_{\min}}\rho_U 4\pi r^2 dr \frac{4k_{\min}}{4\pi r_1 r} e^{-3.56 k_{\min} r}\cos(-k_{\min} r) \\ &\approx \frac{2m_1}{r_1}\frac{dk_{\min}}{\pi}\rho_U 4 r e^{-3.56 k_{\min} r}\cos(-k_{\min} r) dr\end{aligned} \qquad (5.2.\ 4)$$

This is the virtual graviton density at point P due to each spherical shell. The total graviton density is (putting $k = k_{\min}$, $\Upsilon = k_{\min} R_{CCH}$ from Eq. (5.1. 8) then integrating over radius r)

$$\begin{aligned}\Delta\rho_{Gk\min} &= \frac{2m_1}{r_1}\frac{dk_{\min}}{\pi} 2\rho_U \int_0^{R_{CCH}} 2re^{-3.56 k_{\min} r}\cos(-k_{\min} r) dr \\ &= \frac{m_1}{r_1}\frac{dk_{\min}}{k_{\min}^2}\frac{4\rho_U}{\pi}\int_0^{\Upsilon} 2(k_{\min} r) e^{-3.56(k_{\min} r)}\cos(-k_{\min} r) d(k_{\min} r)\end{aligned} \qquad (5.2.\ 5)$$

Thus the extra virtual graviton density $\Delta\rho_{Gk\min}$ at point P distance r_1 from mass m_1 is

Again for brevity putting $\int_0^{\Upsilon} 2(k_{\min}r)e^{-3.56(k_{\min}r)}\cos(-k_{\min}r)d(k_{\min}r) = B$

$$\Delta\rho_{Gk\min} = \frac{m_1}{r_1}\frac{dk_{\min}}{k_{\min}{}^2}\frac{4\rho_U}{\pi}\cdot B \qquad (5.2.\ 6)$$

But we have conjectured that the expansion of space due to GR is such that extra gravitons near a local mass concentration do not change the background density $\rho_{Gk_{\min}}$ in Eq. (5.1. 9) . If the local expansion of space due to GR is $\Delta V/V$ and the new background graviton density $\rho_{Gk_{\min}}{}'$ is to remain unchanged then

$$\text{New } \rho_{Gk\min}{}' \approx \frac{\rho_{Gk_{\min}} + \Delta\rho_{Gk\min}}{1+\Delta V/V} \approx \text{ original } \rho_{Gk\min} \text{ implying } \frac{\Delta\rho_{Gk\min}}{\rho_{Gk\min}} \approx \frac{\Delta V}{V} \qquad (5.2.\ 7)$$

Thus using Eq's. (5.1. 9) & (5.2. 6)

$$\frac{\Delta V}{V} \approx \frac{\Delta\rho_{Gk\min}}{\rho_{Gk\min}} \approx \frac{\left[\dfrac{m_1}{r_1}\right]\dfrac{dk_{\min}}{k_{\min}{}^2}\dfrac{4\rho_U}{\pi}\cdot B}{\dfrac{dk_{\min}8\rho_U{}^2}{k_{\min}{}^4}\cdot A} \approx \left[\frac{m_1}{r_1}\right]\left[\frac{k_{\min}{}^2}{2\pi\rho_U}\frac{B}{A}\right] \qquad (5.2.\ 8)$$

Both integrals B & A are functions of $\Upsilon = k_{\min}R_{CCH}$. In comoving coordinates, at any cosmic time T, the red part of this equation must be a fixed value at all points inside the horizon and thus the expansion of space around any mass is proportional to m_1/r_1. The Schwarzschild solution to Einstein's equations tells us that the radial metric around mass m_1 changes as

$\dfrac{\Delta r_{local}}{\Delta r_{\infty}} = (1-\dfrac{2m_1}{r_1})^{-1/2} \approx 1+\dfrac{m_1}{r_1}$ when $r_1 >>> m_1$ and the local change in volume $\Delta V/V \approx m_1/r_1$.

We have been approximating to the first order in m_1/r_1 so to this first order we can say

$$\frac{\Delta V}{V} \approx \frac{m_1}{r_1} \approx \left[\frac{m_1}{r_1}\right]\left[\frac{k_{\min}{}^2}{2\pi\rho_U}\frac{B}{A}\right] \qquad (5.2.\ 9)$$

Provided the background graviton density $\rho_{Gk\min}$ remains constant, GR tells us that, if working in Planck units, the red highlighted part is approximately one and so

The average density of the universe $\rho_U \approx \frac{k_{\min}{}^2}{2\pi}\frac{B}{A}$ (5.2. 10)

We can now put this value of the average density into Eq.(5.1. 9) for $\rho_{Gk\min}$.

$$\text{Graviton density @ } k_{\min} = \rho_{Gk\min} = \frac{8dk_{\min}}{k_{\min}{}^4}\rho_U{}^2 \cdot A = \frac{8dk_{\min}}{k_{\min}{}^4}\left[\frac{k_{\min}{}^4}{4\pi^2}\frac{B^2}{A^2}\right] \cdot A = \frac{2dk_{\min}}{\pi^2}\frac{B^2}{A} \quad (5.2.\ 11)$$

If our conjecture is true, this is the average density of gravitons at wavenumber $k_{\min}$ excluding possible effects of virtual particles emerging from the vacuum. In section 6.2.1 we will argue these do not change Eq. (5.2. 11) and look next at the metric near large masses.

5.2.2 The Schwarzchild metric near large masses.

At a radius r_1 from a mass m_1 the Schwarzchild metric is $(1-2m_1/r_1)^{\pm 1/2}$ for the time and radial terms. The radial term can be written as

$$\frac{1}{\sqrt{1-2m_1/r_1}} = \frac{1}{\sqrt{1-\beta^2}} = \gamma \quad (5.2.\ 12)$$

Velocity β $(c=1)$ is that reached by a small mass falling from inifinity and $\gamma^{\pm 1}$ is the metric change in clocks and rulers due to mass m_1. Differentiating the metric at fixed radius r_1

$$d\left[1-\frac{2m_1}{r_1}\right]^{-1/2} = \frac{dm_1}{r_1}\left[1-\frac{2m_1}{r_1}\right]^{-3/2} = \gamma^3\frac{dm_1}{r_1}$$

We can write this as the change in the radial metric $\Delta\gamma = \gamma^3\frac{\Delta m_1}{r_1}$ (5.2. 13)

Where $\Delta\gamma$ is the change in the metric from adding mass Δm_1 to mass m_1. With zero mass $m_1 = 0$ the unexpanded unit volume $V = 1$. When $m_1 > 0$ the volume $V = \gamma$ so that

$$\frac{\Delta V}{V} = \frac{\Delta\gamma}{\gamma} = \gamma^2\frac{\Delta m_1}{r_1} \quad (5.2.\ 14)$$

Starting with a mass m_1 bring in from infinity a small extra mass Δm_1 and repeat the derivation of Eq. (5.2. 6) in section 5.2.1. Because of the metric γ at point P in Figure 5.2. 1

due to m_1 compared to infinity, an observer measures the total mass of all the spherical shells inside his causally connected horizon as $\gamma\int(dm)$. Also the mass m_1 in the Schwarzchild metric is that of the mass dispersed at infinity before it comes together. Thus an observer at P also measures the added mass Δm_1 as $\gamma(\Delta m_1)$. These two factors of γ together modify Eq. (5.2. 6) where $\Delta\rho_{Gk\min}$ is now the extra gravitons at point P from adding mass Δm_1 :

$$\Delta\rho_{Gk\min} \approx \gamma^2\left[\frac{\Delta m_1}{r_1}\right]\frac{dk_{\min}}{k_{\min}{}^2}\frac{4\rho_U}{\pi}\cdot B \tag{5.2. 15}$$

Equation (5.2. 9) is also modified, where ΔV is now the volume expansion at point P predicted by GR from adding mass Δm_1 : $\dfrac{\Delta V}{V}\approx\gamma^2\dfrac{\Delta m_1}{r_1}\approx\gamma^2\left[\dfrac{\Delta m_1}{r_1}\right]\left[\dfrac{k_{\min}{}^2}{2\pi\rho_U}\dfrac{B}{A}\right]$

Again the part in red is equal to approximately one and we have agreement with Eq. (5.2. 14)

$$\frac{\Delta V}{V}\approx\frac{\Delta\rho_{Gk\min}}{\rho_{Gk\min}}\approx\frac{\Delta\gamma}{\gamma}\approx\gamma^2\frac{\Delta m_1}{r_1} \tag{5.2. 16}$$

Thus the Schwarzchild metric around any mass concentration can be consistent with the conjecture that the background graviton density $\rho_{Gk\min}$ is uniform at any cosmic time T for comoving observers. But we have not yet discussed the effect of the metric on the local measured value of $k'_{\min}=\gamma k_{\min}$. If we look at our derivations again in the above, the key factor in our integrals is the product $k_{\min}r$; in particular $e^{-3.56k_{\min}r}$ & $\cos(-k_{\min}r)$. The wavelength $\lambda_{\max}=1/k_{\min}$ spans approximately to the horizon. There is an extremely localized perturbation of the exponentially decaying wavefunction. Even close to large black holes in the region where the metric $\gamma>>1$ the product $k'_{\min}r\to\gamma k_{\min}r\to 0$ and the overall effect on the integrals is insignificant. The metric γ rapidly $\to 1$ with radius r, and $\gamma k_{\min}$ also rapidly $\to k_{\min}$. We also use the fact that outside observers see infalling mass remaining on the event horizon, and gravitons emitted from the horizon as seen by this observer. Also we assumed in our derivations that space is approximately flat. For the same reasons as above, as the region around even large black holes (where space is far from flat) is insignificant in relation to the volume over which we integrate, the overall effect is again very small.

5.3 The Expanding Universe

The big problem with the above sections is that the density of zero point modes available at $k_{\min} \approx 1/R_{CCH}$ is too small by a factor of approximately the area of the causally connected horizon R_{CCH}. The value we calculated for $\rho_{Gk\min}$ in Eq. (5.2. 11) is $\approx R_{CCH}{}^2$ times greater than the zero point modes available. Each graviton of wavenumber $k_{\min} \approx 1/R_{CCH}$ (as in section 6.2) has rest mass $m_0 \approx 1/R_{CCH}$ with a similar momentum (with $c=1$), so the velocity of its rest frame is such that the density of zero point modes required to build these graviton superpositions is also approximately $\approx \rho_{Gk\min}$. The density of zero point modes available at $k_{\min} \approx 1/R_{CCH}$ however is $\dfrac{k_{\min}{}^2}{\pi^2} \approx \dfrac{1}{\pi^2 R_{CCH}{}^2}$, so how can this be?

5.3.1 Holographic horizons and red shifted Planck scale zero point modes.

Zero Point energy transformations are invariant in all rest frames. The Lorenz transformations tell us this, but it is due to Special Relativity which applies locally. In section 2.2.3 we defined a rest frame in which zero momentum preons with infinite wavelength build infinite superpositions. If we have a spherical horizon of space with Planck scale modes expanding at the velocity of light, these Planck modes can be absorbed by infinite wavelength preons and red shifted in a radially focussed manner inwards. We will argue in what follows, that at the centre (where the infinite superpositions are built), approximately 1/3 of the Planck modes on the horizon become available with wavelengths of the order of the radius of this horizon. We are also assuming (as in sections above) that GR applies at cosmic scales but not all the way to the horizon. We also temporarily ignore the fact that the *Hubble horizon* is seen as the radius at which space expands at the speed of light and replace it with the *causally connected horizon*. We return to these points in sections 5.3.2 & 5.3.3

Starting with Eq. (5.1. 8) $\Upsilon = k_{\min} R_{CCH}$ let space at the horizon expand radially away from a central observer at velocity $c=1$. The Planck scale modes will be red shifted to wavelengths $R_{CCH} \approx 1/k_{\min}$ etc. Let $r=$ the radial distance inwards from the horizon R_{CCH}. In Planck units:

The space expansion velocity $$\beta = \frac{v}{c} = 1 - \frac{r}{R_{CCH}} \to 1 \text{ when } r = 0 \text{ @ } R_{CCH}. \qquad (5.3.\ 1)$$

Using Eq. (5.3.1) define $$k' = \frac{k_{\min}}{1-\beta} = \frac{k_{\min} R_{CCH}}{r} = \frac{\Upsilon}{r} \qquad (5.3.\ 2)$$

Where k' is the wavenumber @ *r* that redshifts to $k_{\min}$. Also $k' = \Upsilon / r = 1$ when $r_{\text{minimum}} = \Upsilon$.

This is the Planck region.

Using Eq's. (5.3. 1) & (5.3. 2) $$dk' = \frac{dk_{\min} R_{CCH}}{r} - \frac{k_{\min} R_{CCH} dr}{r^2}$$
$$= \frac{dk_{\min}}{k_{\min}} \frac{k_{\min} R_{CCH}}{r} - k_{\min} R_{CCH} \frac{dr}{r^2}$$

$$dk' = \frac{\Upsilon}{r} \frac{dk_{\min}}{k_{\min}} - \Upsilon \frac{dr}{r^2} \tag{5.3. 3}$$

The density of modes @ $k' \ = \dfrac{k'^2 dk'}{\pi^2}$. In a layer of the shell thickness dr and unit area, the number of modes using Eq. (5.3. 2) is $\dfrac{dN}{A} = \dfrac{k'^2 dk' dr}{\pi^2} = \dfrac{\Upsilon^2}{\pi^2 r^2} \left[\dfrac{\Upsilon dr}{r} \dfrac{dk_{\min}}{k_{\min}} - \dfrac{\Upsilon (dr)^2}{r^2} \right]$ where in the limit the second term can be ignored and thus

$$\frac{dN}{A} = \frac{k'^2 dk' dr}{\pi^2} = \frac{\Upsilon^3}{\pi^2} \frac{dk_{\min}}{k_{\min}} \frac{dr}{r^3}$$

$$\frac{N}{A} = \frac{\Upsilon^3}{\pi^2} \frac{dk_{\min}}{k_{\min}} \int_{r(\min)}^{R_{CCH}} \frac{dr}{r^3} \approx \frac{\Upsilon^3}{\pi^2} \frac{dk_{\min}}{k_{\min}} \frac{1}{2 r_{\min}{}^2} \tag{5.3. 4}$$

Only 1/3rd of these modes are transverse to radial directions and available to redshift radially in a focussed manner to build infinite superpositions. We also assume the smallest value of $r_{\min}$ is the Planck length with $r_{\min} = 1$. The surface area *A* of the horizon is $A = 4\pi R_{CCH}{}^2$.

The total number of modes available at $k_{\min}$ using Eq. (5.3. 4) is:

$$N_{Total} = 4\pi R_{CCH}{}^2 \times \frac{1}{3} \times \frac{N}{A} \approx 4\pi R_{CCH}{}^2 \frac{1}{3} \frac{\Upsilon^3}{\pi^2} \frac{dk_{\min}}{k_{\min}} \frac{1}{2}$$

As the wavelengths of these modes are $\lambda \approx R_{CCH}$ they occupy the volume inside the horizon with effective density $\rho_{k\min} \approx \dfrac{N_{Total}}{4\pi R_{CCH}{}^3 / 3} \approx \dfrac{4\pi R_{CCH}{}^2}{(4\pi R_{CCH}{}^3 / 3)} \dfrac{\Upsilon^3}{\pi^2} \dfrac{dk_{\min}}{k_{\min}} \dfrac{1}{6} \approx \dfrac{3}{R_{CCH}} \dfrac{\Upsilon^3}{\pi^2} \dfrac{dk_{\min}}{k_{\min}} \dfrac{1}{6}$

But $\Upsilon = k_{\min} R_{CCH}$ and so $$\rho_{k\min} \approx \frac{\Upsilon^2}{2\pi^2} dk_{\min} \tag{5.3. 5}$$

There is another factor however. The analysis above assumes that the horizon is at a fixed radius with space expanding through it. As the horizon is also expanding this increases the effective thickness of the shell. The zero point Planck modes on the shell using Eq.(5.3. 2) last for time $\Delta t \approx \hbar / \Delta E \approx 1/k'$. For example at the horizon where $r_{\min} = 1$ and $k' = 1$ then $\Delta t \approx 1$. If the horizon is moving at $c = 1$ say, in the time $\Delta t \approx 1$ it travels a distance $\Delta x = \Delta t \approx 1$ and the effective thickness of the horizon is doubled. If we call this velocity of the horizon *V*, the effective thickness increases as $dr' = (V+1)dr$ and Eq. (5.3. 5) becomes

$$\rho'_{k\min} \text{ available } = (V+1)\rho_{k\min} = (V+1)\frac{\Upsilon^2}{2\pi^2}dk_{\min} \tag{5.3. 6}$$

Equation (5.2. 11) is the density of gravitons at $k_{\min}$ as a function of Υ and each of these gravitons has to absorb a zero point mode of the same wavenumber. We can thus say:

$$\rho_{k\min} \text{ required } = \frac{2}{\pi^2}\frac{B^2}{A}dk_{\min} \tag{5.3. 7}$$

Figure 5.3. 1 superimposes plots of Eq. (5.3. 6) as the available $k_{\min}$ modes over a wide range of possible velocities $V = 1 \rightarrow 6$ heading towards the future in an accelerating manner. It also plots the required $k_{\min}$ modes in Eq. (5.3. 7), all as functions of $\Upsilon = k_{\min}R_{CCH}$.

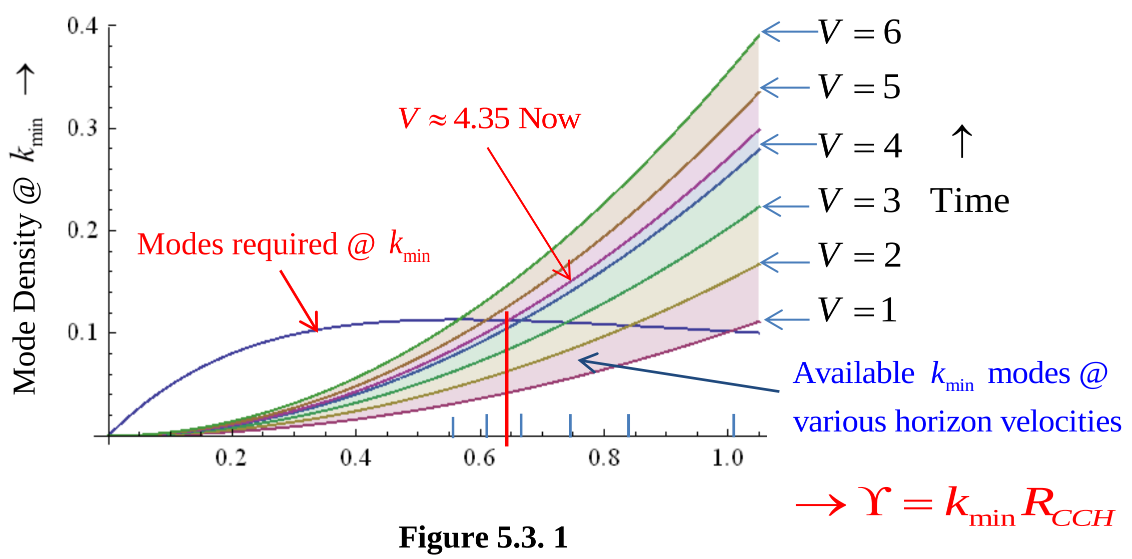


Figure 5.3. 1

Figure 5.3. 1 tells us the values of $\Upsilon = k_{\min}R_H$ where the zero point modes available $@k_{\min}$ equal the zero point modes required for various receding horizon velocities. We can then put these values of Υ into Eq. (5.2. 10) for the average density of the universe, but first rewriting it as Eq. (5.3. 8) which is plotted in Figure 5.3. 2.

$$\rho_U \approx \frac{k_{\min}{}^2}{2\pi}\frac{B}{A} = \frac{\Upsilon^2}{2\pi R_{CCH}{}^2}\frac{B}{A} \qquad \text{or} \qquad \rho_U R_{CCH}{}^2 \approx \frac{\Upsilon^2}{2\pi}\frac{B}{A} \tag{5.3. 8}$$

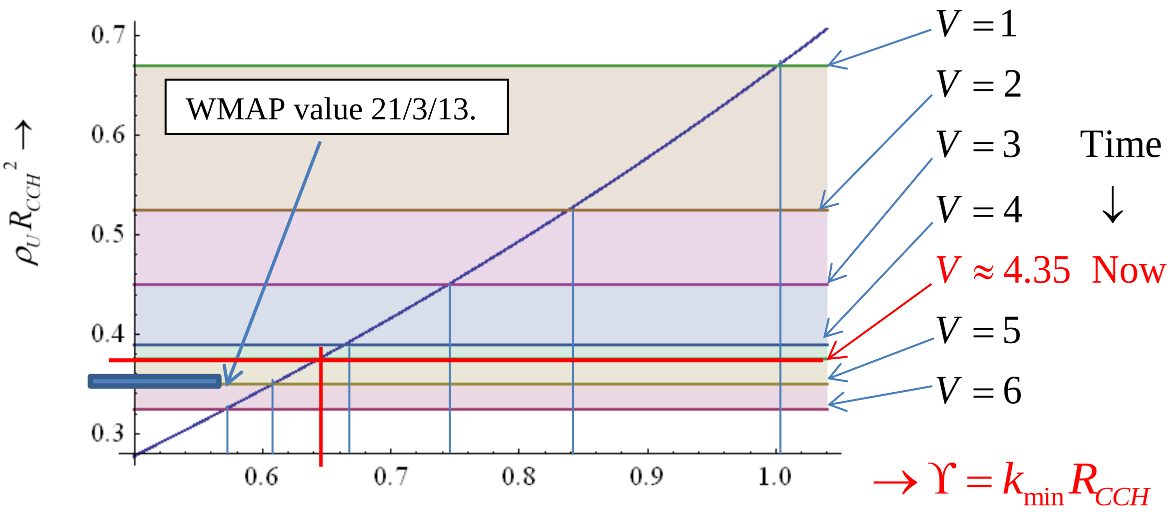


Figure 5.3. 2

Figure 5.3. 2 shows at constant $\Upsilon = k_{\min} R_{CCH}$ the density of the universe decreases $\propto 1/R_{CCH}{}^2$. If initially Υ starts at ≈ 1 say, and steadily decreases for some reason, the density decrease is $> 1/R_{CCH}{}^2$ and could appear as acceleration. Also at a horizon velocity of $V \approx 5$ the baryonic plus dark matter density of the universe is the same as the 9 year WMAP data (21 March 2013). The FLRW models, which are based on GR reaching to the horizon boundary (*we have been arguing in the above that it doesn't, but more on that below*), tell us that in the very early eras the mass energy density is radiation dominated, the scale factor varies as $a \propto t^{1/2}$ and the (causally connected) horizon velocity (see Eq's. (5.3. 10) is $V \rightarrow 2(c)$. In the later matter dominated eras $a \propto t^{2/3}$ and $V \rightarrow 3$. The red shift data tell us the Hubble constant is currently $H \approx 1/t$, the scale factor is approximately $\propto t$ with a current horizon velocity of $V \approx 4.35$ (see Eq's. (5.3. 10). This acceleration is widely attributed to dark energy, which the WMAP values assume in making the required total energy density of $\Omega = 1$ (again assuming GR reaches to the horizon). Also if the current horizon velocity is $V \approx 4.35$ Figure 5.3. 2 shows a matter density $\approx 7\% >$ than WMAP. The 9 year WMAP values are based on a Hubble constant of 69.3 and at ≈ 72 it would match the WMAP values. The WMAP dark matter values are estimated assuming dark energy is required to make the total energy density $\Omega = 1$, but if this is not needed then slightly higher dark matter values could still match the WMAP data.

5.3.2 General Relativity may not reach to the Horizon, and the implications

It is nearly 100 years since Einstein published his General Relativity. It has also always been assumed that the FLRW metric, which is an exact solution to Einstein's field equations, applies all the way to the boundaries of any causally connected universe. In section 5.2.1 however the approximations we made meant that the agreement we get with GR in Eq. (5.2. 14) is true only for points close to a mass concentration in relation to the distance to the horizon R_{CCH}. This implies that GR or in turn the expansion of space around any mass concentration is locally true, and true to some substantial fraction of R_{CCH}, probably some billions of light years, but not true at distances of the order of R_{CCH}. Also this agreement is true only when we assumed that space is essentially flat on average. This does not seem unreasonable if in fact GR does not reach to the horizon; there would then be no effective gravitational field acting at these vast distances trying to distort space on average. Inflation was originally proposed to solve what has always been seen as the overall average flatness problem, but it may in fact not be necessary if the above is true. Of course in the initial very early era, before the causally connected region was large enough to allow infinite superpositions, some form of inflation is not ruled out by the above. Summarizing all this:

- *Local warping of spacetime depends only on local mass concentrations*.
- *The causally connected universe is flat on average* (Ω *does not need to be one for this to be true).*
- *The overall average expansion of space depends on Eq.*(5.3. 8).

5.3.3 Scale factors and the expansion of space.

In section 5.3.1 we made the assumption that space is expanding at the velocity of light at the causally connected horizon and not the Hubble horizon, how can we possibly justify this? The stretching of the scale factor between comoving distant galaxies has always been seen as equivalent to the expansion of space between them and is an integral part of the FLRW solutions to Einstein's field equations. But again this assumes that GR reaches the causally connected horizon as in section 5.3.2. So if GR does not reach the horizon it does not determine the overall expansion of space as in our final paragraph of section 5.3.2. *If Eq.* ((5.3. 8**)** *determines the average density of the universe it also determines the spacing between distant galaxies and hence what has always been referred to as the scale factor*. If this is so the expansion between distant galaxies does not require local forces. If all comoving galaxies measure the same cosmic time T after the big bang they must be moving on local geodesics. If there were any local forces operating, the elapsed time would not be the maximized time that occurs along geodesic paths, because GR is working locally. Thus we argue; *there does not need to be conflict between differing overall expansion rates for space and galaxies.* The same applies to any waves in space but there are some differences. As a simplified example imagine uniformly distributed continuous standing waves spanning all space. Between the nodes of these standing waves are energy or mass concentrations also moving on geodesic paths, just as the distant galaxies do. The distance between nodes will stretch with time as if they are on comoving coordinates. *This red shift with time depends only on the scale factor (between galaxies) exactly as in current cosmology.* But what effect does the expansion of space itself have, as we are arguing that it is different to the expansion between galaxies? In section 5.3.1 we used the infinite wavelength of zero momentum preons (section 2.2.3) to borrow redshifted Planck modes from a holographic horizon of space expanding at the velocity of light. This is *instantaneous redshift acting across vast distances, of a quantum mechanical nature, and due to the expansion of space itself.* The redshift we observe from distant galaxies is due to the scale factor changing over vast periods of time. *One is time related; the other instantaneous, quantum mechanical, and distance related.* The latter is the same instantaneous connection over vast distances that so deeply troubled Einstein. Looking at scale factors and space expansion a little further; if the scale factor is $a(t) \propto t^p$ with $\dot{a}(t) \propto pt^{p-1}$ the Hubble parameter $H(t) = \dfrac{\dot{a}(t)}{a(t)} = \dfrac{p}{t}$, and at the present time T

Distance to the Hubble horizon using $H(T)R_{Hubble} = c = 1$ is $R_{Hubble} = \dfrac{1}{H(t)} = \dfrac{T}{p}$ (5.3. 9)

The Hubble horizon radius R_{Hubble} is linear with time *T*, but inverse with *p*, and this is always true. On the other hand the causally connected horizon radius R_{CCH}, and its velocity *V*, depend on the behaviour of *p* if it varies as a function of time. Writing the present scale factor, normalized to one, as $a(T)$ and $a(t) \propto t^p$, then using Eq. (5.3. 9) it follows that:

The horizon radius $R_{CCH} = a(T)\displaystyle\int_0^T \frac{dt}{a(t)} = T^p \int_0^T \frac{dt}{t^p} = \frac{T}{1-p}$ only when p is constant.

The horizon velocity $V = \dfrac{1}{1-p}$ only when p is constant. (5.3. 10)

But horizon velocity $V = \dfrac{dR_{CCH}}{dT} = 1 + H(T)R_{CCH} = 1 + \dfrac{p}{T}R_{CCH}$ is always true.

From these equations as $R_{CCH}/T \approx 3.35$ and $p \approx 1$ at present, the horizon velocity $V \approx 4.35$. If our arguments are correct, the expansion of space equals the expansion of distant galaxies (or comoving coordinates) only when $R_{Hubble} = R_{CCH}$ at a constant $p = 1/2$, or the constant horizon velocity $V = 2$. This is only true in the radiation dominated phase of current cosmology, which shows identical behaviour to our conjecture above. Even though the minimum possible horizon velocity has to be at least $V = 1$, it appears that $V = 2$ must be the starting point in both cases; all nucleosynthesis calculations in this early phase of expansion assume that $\rho_U \propto 1/t^4$ or equivalently $p = 1/2$. The interesting point however is that if the horizon velocity steadily increases, the horizon radius is proportional to the average velocity, but the modes available in Eq. (5.3. 6) depend on the peak velocity *V*. There may be some sort of positive feedback here; any increase in *V* allows density changes in Eq. (5.3. 8), which possibly becomes an ongoing process leading to the $V \approx 4.35$ we now see. Possibly *V* has steadily increased with time from its starting point of $V = 2$ to the current value. Of course the rate of change with time has to be in agreement with the redshifts currently observed when looking back towards the big bang. Also recent observations have been questioning the leading current dark energy explanations of acceleration [1].

6 Further consequences of Infinite Superpositions

6.1 Low frequency Infinite Superposition cutoffs

In section 4.2 when we introduced gravity, for the lower limit in our integrals we assumed $k_{\min} = 0$, and then in section 5 showed that there is a lower limit $k_{\min} > 0$. It turns out that for massive $N = 1$ superpositions the effect of this is negligible in comparison to the high frequency cutoff $k_{cutoff} < \infty$, which we showed gravity can address in section 4.2. For infinitesimal rest mass $N = 2$ superpositions we cannot however ignore the effect of $k_{\min} > 0$.

6.1.1 Quantifying the approximate effect of $k_{\min} > 0$ on infinite superpositions

If we look again at section 4.2.1 we can repeat what we did there as follows. Initially to illustrate these effects we will consider only $N = 1$ superpositions where we can say that

$$\text{When } K_{nkCutoff} \to \infty \ \ \& \ \ (\text{for } N = 1 \text{ only}) \ K_{nk\min} \to 0 \tag{6.1. 1}$$

$$\left[\frac{-1}{1+K_{nk}{}^2}\right]_{K_{nk\min}}^{K_{nk}cutoff} = \frac{1}{1+K_{nk\min}{}^2} - \frac{1}{1+K_{nkCutoff}{}^2} \approx 1 - \left[\frac{1}{K_{nkCutoff}{}^2} + K_{nk\min}{}^2\right] \approx \frac{1}{1+\varepsilon'}$$

Our infinitesimal $\varepsilon \to \varepsilon' \approx \dfrac{1}{K_{nkCutoff}{}^2} + K_{nk\min}{}^2$ and from Eq. (3.1. 11) $K_{nk}{}^2 = \dfrac{n^2 s}{2} ƛ_c{}^2 k^2$

For spin ½ fermions for example $\langle n^2 s/2 \rangle \approx 9$. Also $k_{cutoff}{}^2 \approx 1/L_P{}^2$ and $k_{\min}{}^2 \approx 1/R_H{}^2$ so that

$$\varepsilon' \approx \frac{1}{K_{nkCutoff}{}^2} + K_{nk\min}{}^2 \approx \frac{L_P{}^2}{9ƛ_c{}^2} + \frac{9ƛ_c{}^2}{R_{CCH}{}^2} \approx \frac{(L_P R_{CCH})^2 + (9ƛ_c{}^2)^2}{9ƛ_c{}^2 R_{CCH}{}^2}$$

The ratio of the extra contribution $\Delta\varepsilon$ to ε where $\varepsilon' = \varepsilon + \Delta\varepsilon$ is $\dfrac{\Delta\varepsilon}{\varepsilon} \approx \left[\dfrac{9ƛ_c{}^2}{L_P R_{CCH}}\right]^2$ (6.1. 2)

Equation (6.1. 2) is for spin ½ but the numerical factor 9 only changes slightly for spins 1 & 2. Working in Planck units $L_P R_{CCH} \approx 10^{61}$, but for electrons say $ƛ_c{}^2 \approx 6\times10^{44}$, so the effect is of the order of $\Delta\varepsilon/\varepsilon \approx 10^{-30}$ which we have been ignoring. But we cannot ignore this in the case of infinitesimal rest masses as we will see.

6.2 Infinitesimal Masses and *N* = 2 Superpositions

Looking again at angular momentum and rest masses in section 3.2 the key factor in our final integrals is in Eq. (6.1. 1). Using Eq. (3.1. 12) we can rewrite Eq. (6.1. 1) as

$$\left[\frac{-1}{1+K_{nk}{}^2}\right]_{K_{nk\min}}^{K_{nk}cutoff} = \frac{1}{\gamma_{nk\min}{}^2} - \frac{1}{\gamma_{nk\mathrm{Cutoff}}{}^2} \tag{6.2. 1}$$

With massive $N=1$ superpositions as above the difference between $\gamma^2{}_{nk\min}$ & 1 is vanishingly small, i.e. $(\gamma^2{}_{nk\min} - 1) \to 1/\infty$ and as in section 6.1.1 this first term is of much less significance than the $\gamma^2{}_{nk\mathrm{cutoff}}$ term. But now define the *N* of Eq's. (2.1. 4) & (2.2. 4) using Eq. (3.1. 12) as follows

$$N = 1 + \langle K_{nk\min}{}^2 \rangle = \langle \gamma_{nk\min}{}^2 \rangle \tag{6.2. 2}$$

In section 3.2 we derived angular momentum and rest masses for only massive or what we called $N=1$ particles. To get integral angular momentum we had to assume in deriving Eq. (3.2. 6) that the minimum value of K_{nk} or $K_{nk\min} = 0$. For massive $N=1$ particles such as the fermions the error in this assumption (as in section (6.1. 1) is $\approx 10^{-30}$ times smaller than ε, which for an electron is already $\varepsilon \approx 10^{-45}$ due to the high frequency cutoff $@ \approx 10^{18.31} GeV$. (We allowed for this $\varepsilon \approx 10^{-45}$ when we included gravity in section 4.2.) From section (6.1. 1) above we approximated $K_{nk\min}{}^2$ as $\approx 9\bar{\lambda}_c{}^2 / R_H{}^2$ for a spin ½ fermion. So we can express Eq. (6.2. 2) in terms of this approximation

$$N \approx 1 + \frac{9\bar{\lambda}_C{}^2}{R_H{}^2} \approx 1 \quad \text{as} \quad \frac{9\bar{\lambda}_C{}^2}{R_H{}^2} \to 0 \tag{6.2. 3}$$

$$\text{For example an electron } \frac{9\bar{\lambda}_C{}^2}{R_H{}^2} \approx 10^{-77}$$

For most of the massive particles it appears that $N \to 1$. Even if neutrino masses were as low as $10^{-4} eV$ then $N - 1 \approx 10^{-59}$. If the mass is too small however we cannot get the correct angular momentum unless something else changes. Infinitesimal increases above 1 of the order of $\approx 10^{-50}$ or so can be handled perhaps by a small change in the actual high frequency cutoff details but this probably does not allow massive particles to be much less than sub micro electron volts. So if massive particles are a group with $N \approx 1$, then it would not seem unreasonable to imagine there could possibly be another group with $N = 1 + \langle K_{nk\min}{}^2 \rangle \approx 2$ implying that $\langle K_{nk\min}{}^2 \rangle \approx 1$. Repeating the derivation of Eq. (3.2. 6) but with $N = 1 + \langle K_{nk\min}{}^2 \rangle \to 2$ and for clarity and simplicity let $K_{nk\mathrm{cutoff}} \to \infty$.

$$\mathbf{L}_z(Total) = s \cdot (N = 2) m\hbar \int_{K_{nk\min}}^{\infty} \frac{K_{nk}{}^2}{(1 + K_{nk}{}^2)^2} \frac{dK_{nk}}{K_{nk}} = sm\hbar \left[\frac{-1}{1 + K_{nk}{}^2} \right]_{K_{nk\min}}^{\infty} \tag{6.2. 4}$$

$$\mathbf{L}_z(Total) = sm\hbar \left[\frac{1}{1 + K_{nk\min}{}^2} \right] = sm\hbar \left[\frac{1}{(N = 2)} \right] = \frac{sm\hbar}{2} \text{ as previously.}$$

Provided we have doubled the probability of superpositions as in Eq. (2.1. 4) from $s \cdot (N = 1) dk / k$ to $s \cdot (N = 2) dk / k$, the final angular momentum results in Eq's. (3.2. 6) & (6.2. 4) are identical. The same is true for rest mass calculations. For complete infinite superpositions if $N = 2$ then the expectation value $\langle K_{nk\min}{}^2 \rangle = 1$. We thus conjecture that all $N = 2$ infinite superpositions have $\langle K_{nk\min}{}^2 \rangle \approx 1$.

From Table 4.3. 1; $N = 2$ infinitesimal rest mass spin 1 superpositions have $\langle n^2 \rangle \approx 16.77$
$N = 2$ infinitesimal rest mass spin 2 superpositions have $\langle n^2 \rangle \approx 11.69$

Using Eq's. (5.1. 8) and (3.1. 12)

$$\left\langle K_{nk\min}{}^2 \right\rangle = \frac{\left\langle n^2 \right\rangle s}{2} \lambda\!\!\!^-{}_c{}^2 k_{\min}{}^2 \approx \frac{16.77}{2} \lambda\!\!\!^-{}_C{}^2 k_{\min}{}^2 \approx 1 \text{ or } \lambda\!\!\!^-{}_C \approx 0.345 \frac{R_{CCH}}{\Upsilon} \text{ for Spin 1}$$
$$\approx \frac{11.69 \times 2}{2} \lambda\!\!\!^-{}_C{}^2 k_{\min}{}^2 \approx 1 \text{ or } \lambda\!\!\!^-{}_C \approx 0.292 \frac{R_{CCH}}{\Upsilon} \text{ for Spin 2} \tag{6.2. 5}$$

At the current horizon velocity of $V \approx 4.35$ Figure 5.3. 1 tells us that $\Upsilon \approx 0.646$ and at the current horizon radius of $\approx 46 \times 10^9$ light years $R_{CCH} \approx 2.7 \times 10^{61}$ Planck lengths.

Spin	Compton Wavelength $\lambda\!\!\!^-{}_C$	Infinitesimal Rest Mass
1	$\approx 0.534 R_{CCH}$	$\approx 2.33 \times 10^{-34}$ eV.
2	$\approx 0.452 R_{CCH}$	$\approx 1.97 \times 10^{-34}$ eV.

Table 6.2 1 Infinitesimal rest masses of $N = 2$ photons, gluons & gravitons.

These Compton wavelengths and rest masses are the present values, slowly changing with cosmic time *T.* Also both the electromagnetic force and gravity reach about half way to the horizon. The graviton rest masses above are close to recent proposals explaining the accelerating expansion of the cosmos [2] [3]. As $N = 1 + \left\langle K_{nk\min}{}^2 \right\rangle = \left\langle \gamma_{nk\min}{}^2 \right\rangle$ all $N = 2$ superpositions have $\left\langle \gamma_{nk\min}{}^2 \right\rangle = 2$. Table 6.2 1 however may not be the only $N = 2$ possibilities, but they do appear to be the only ones cutting off at $k_{\min} \approx 1/R_{CCH}$ and we have called them infinitesimal rest masses. *However there seems to be no reason stopping more massive* $N = 2$ *versions with much more energetic low frequency cutoffs,* $k_{\min} \approx m_0$ (i.e. approximately equal to their rest mass) provided $1 + \left\langle K_{nk\min}{}^2 \right\rangle = \left\langle \gamma_{nk\min}{}^2 \right\rangle = 2$. These may have some interesting consequences in relation to massive particles emerging from the vacuum.

6.2.1 Massive *N* = 2 virtual particles emerging from the vacuum?

Is it possible that all the massive virtual particles emerging from the vacuum are $N = 2$ superpositions? If their infinite superpositions cutoff at a minimum wavenumber $k_{\min} > 1/R_{CCH}$ they cannot couple (see section 3.3) to wavelengths of the order of the horizon radius, and this must effect how they relate to GR. Virtual particles of rest mass m_0 only last for time $\Delta t \le 1/m_0$ in natural Planck units so if their superpositions cutoff at $k_{\min} \approx m_0$ they would still behave just as a virtual particle with rest mass m_0 should. For almost a century many papers have been written about why enormous zero point energy densities at Planck scale don't appear to cause a huge curvature of space. In section 5.2.1 we argued that the curvature of space is consistent with a constant background density of the lowest possible energy gravitons $\rho_{Gk\min}$ as in Eq. (5.1. 9). We thus argue that if all virtual massive particles are $N=2$ superpositions they will not influence this background $\rho_{Gk\min}$. However there still needs to be a sufficient density of zero point modes available at the higher cutoff frequencies $k_{\min} \approx m_0$ of these massive virtual particles. Planck scale virtual particles are moving with Planck energies. The rest frame in which they are built (section 2.2.3) is moving at a velocity equal to their momentum over their rest mass i.e. $v = \beta = k/(m_0\gamma)$ in Planck units. Consider some frame (a) fixed in a set of comoving coordinates. Let frame (b) move along the positive x axis say at velocity $\beta = k_x/(m_0\gamma)$ Zero point modes are uniform in all rest frames so are identical in both (a) and (b). All the zero point virtual particles in (b) that have that have momentum $k_x = -m_0\beta\gamma$ will be at rest in frame (a), and due to special relativity *the spacing between them along the x axis will increase as* γ and their *linear density will be reduced as* $1/\gamma$. This linear density of modes reduction is true for all other directions with relative velocity $\beta = k/(m_0\gamma)$ between frames (a) & (b), and with virtual particle momentums parallel to this new direction. Thus the *overall density of modes as measured in the virtual particles rest frame is also reduced as* $1/\gamma$. The value of $1/\gamma$ depends on which virtual particle pairs are created, but we argue that the redshifted supply of various wavelengths from an expanding horizon (section 5.3.1), combined with the above density reductions, enable sufficient zero point modes to be available to build them.

6.2.2 Contrasting high energy cutoff behaviour of *N* = 1 & *N* = 2 superpositions

Equation (6.2. 1) can be written for both $N=1$ & $N=2$ superpositions using the results of sections 4.2 & 6.2 as follows

$$\left[\frac{-1}{1+K_{nk}{}^2}\right]_{K_{nk\min}}^{K_{nk}cutoff} = \frac{1}{\gamma_{nk\min}{}^2} - \frac{1}{\gamma_{nk\mathrm{Cutoff}}{}^2} = \frac{1}{2(1+\varepsilon')} \quad \text{when } N=2 \qquad (6.2.\ 6)$$

$$= \frac{1}{1+\varepsilon'} \quad \text{when } N=1$$

(We should be using expectation values, but for simplicity will just imply them.) We have just shown above when $N=2$ that $\left\langle 1/\gamma_{nk\min}{}^2 \right\rangle = 1/2$ but in reality it is Eq. (6.2. 6) that has to be true. This allows flexibility in $N=2$ high frequency cutoff behaviour that does not apply to $N=1$ superpositions. (Section 4.2 shows that $N=1$ cutoff near the Planck length.) For example if infinitesimal rest mass photon $N=2$ superpositions, with mass of $\approx 10^{-34} eV$, have a high energy cut off at the $\approx 91 GeV$ of massive photons, then $1/\gamma_{nkCutoff}{}^2$ is $\approx 10^{-90}$.

Thus Eq.(6.2. 6) implies $1/\gamma_{nk\min}{}^2 \approx (1/2) - 10^{-90}$ which seems quite realistic. Summarizing:

- *N=1 superpositions all have high energy cutoffs near the Planck Energy.*
- *N=2 infinitesimal rest mass superpositions on the other hand can have high energy cutoffs many orders of magnitude less than Planck values, but more massive N=2 versions can still cutoff near the Planck energy.*

6.2.3 Are massive photons *N* = 1 or *N* = 2 and real versus virtual particles?

In Table 4.3. 1 we described massive photons as either $N=1$ or $N=2$ and the above attempts to explain why. If they are $N=1$ they must cutoff at $k_{\min} \approx 1/R_{CCH}$ and thus influence gravity. This suggests they live long enough to enable their mass/energy to be measured and could be real. On the other hand if they are massive and $N=2$ they cutoff at $k_{\min} \approx m_0 >> 1/R_{CCH}$ and don't influence gravity, or alternatively don't live long enough, and could be virtual. Taking this argument further however can be problematic, as whether particles are real or virtual can depend on the frame in which measurements are made. However this usually applies to only long wavelength particles escaping black holes or accelerating frames, and thus of infinitesimal rest mass that must be $N=2$. Infinitesimal rest mass particles whether they are virtual or real are both $N=2$ and both cutoff at $k_{\min} \approx 1/R_{CCH}$. But on the other hand massive $\approx 91GeV$ photons have nuclear wavelengths and can only escape sub nuclear size black holes which are thought unstable. So maybe reference frames deciding whether particles are virtual or real is not an issue with massive $N=1$ or $N=2$ photons after all.

6.3 Black Holes, the Firewall Paradox and possible Spacetime Boundaries

Several recent papers [14] [15] [16] [17] [18] have discussed the BH firewall paradox. In section 5.2.2 we use the fact that outside observers see infalling mass remaining on the horizon. In fact if we look carefully at the analyses in sections 5.2.1 & 5.2.2 we see they strongly suggest that GR cutsoff at the BH horizon; one of the possible firewall paradox implications. The equations we derived do not work inside the horizon. Our argument that comoving observers see a constant density of gravitons at $k_{\min}$ being consistent with GR will not work inside the horizon. As Black holes curve spacetime this paper has to argue that all superpositions must be seen from the outside as cutting off at $k_{\min} \approx 1/R_{CCH}$ and thus be $N=1$ for massive and $N=2$ for infinitesimal mass superpositions as distinct from the virtual pair creation in section 6.2.1. To curve the spacetime outside a Black hole, we argue a firewall is essential to convert the enormous real energy of incoming matter into real Planck energy particles. (See Table 6.2 2.)

Infinite Superposition Type	Black Holes	Vacuum virtual pairs
Massive Spin ½ Fermions	$N=1$ near horizon	$N=2$
Massive Spin 1 Bosons	$N=1$ near horizon	$N=2$
Infinitesimal Mass Spin 1 Bosons	$N=2$ above horizon	$N=2$

Table 6.2 2 A tentative scenario: black holes versus vacuum pair creation?

6.4 Dark matter possibilities

Table 4.3. 1 shows a spin 2, $N=1$ neutral massive graviton type superposition that possibly does not interact via the weak force but only with virtual infinitesimal rest mass $N=2$ graviton superpositions. It may possibly be only detected via graviton interactions.

6.5 Higgs Boson superpositions.

Table 2.2. 1 lists a possible Higgs boson group of spin zero preons. A spin zero Higgs boson, if it is in fact a superposition, would have to be some superposition of infinite superpositions with a total angular momentum vector that sums to zero just as two spin ½ fermions can.

6.6 Constancy of fundamental charge

It has always been fundamental that the electromagnetic charge of protons and electrons is precisely equal and opposite to get a neutral universe. In section 4.2 we showed that the probability of superpositions was $sN \cdot dk(1+\varepsilon)/k$ where the infinitesimal ε is proportional to rest mass squared and thus different for various particles. We used this probability to determine interaction coupling strengths in section 3.3. This suggests that the probability of virtual photon emission is also proportional to the probability $sN \cdot dk(1+\varepsilon)/k$ of each superposition, and would not be precisely equal for electrons and protons due to small variations in ε of the order of $\approx 10^{-45}$ between electrons and quarks. If however we look closely again at section 4.2 and Eq. (4.2. 3), by adding the amplitude for gravity at right angles we effectively added the probabilities of spin 2 gravity generated superpositions to those of spin 1 colour and electromagnetic superpositions. If somehow only those superpositions generated by spin 1 electromagnetic and colour interact with spin 1 photons this would cancel any minute difference in charge. If this is not the case then there are infinitesimal differences in charge of the order of $\approx 10^{-45}$ which would probably have shown up in some form unless there are minute differences in the total number of electrons and protons.

6.7 Feynman's Strings

Over a century ago there were various models of the electron. The Abraham-Lorenz was probably the most well-known [19] [20]. All these models suffered from the problem that the electromagnetic mass in the field was 4/3 times the relativistic mass. In 1906 Poincare showed that if the bursting forces due to charge were balanced by stresses (or forces) in the same rest frame as the particle, that these would cancel the extra 1/3 figure, thus restoring covariance [21]. In chapter 29 Volume II of his famous lectures on physics, Feynman, probably jokingly, suggested that if the electron is held together by strings that their resonances could explain the muon mass; he just may have been right [22]. The equations for infinite superpositions in this paper apply equally to all massive particles. Also, as infinite superpositions are held together by interactions with zero point forces in the same rest frame, could these zero point interactions possibly be Feynman's strings? If they hold the virtual preons in orbit, they must also surely be able to balance any bursting forces due to electric charge. However this paper suffers from the same problem as the Standard Model: There is nothing in it suggesting the quantization of mass of massive particles; it suggests only the mass of infinitesimal rest mass particles.

7 Conclusions

If fundamental particles are built from infinite superpositions then why do we never see any sign of them? It is important to remember that all members of infinite superpositions are virtual and only complete infinite superpositions can behave as real particles. If infinite superpositions could be somehow decomposed into their virtual components this would destroy the resulting equivalent real particle. Could it be that particle conservation laws controlling the behaviour of fundamental particles somehow prevent any sign of their virtual components?

Also the viability of this paper depends on primary interactions where spin zero preons can borrow mass from some Higgs type scalar zero point field, and energy from colour and electromagnetic zero point vector fields. The behaviour of these primary interactions is very different to the secondary interactions that the SM is all about. The SM rules applying to borrowing mass and energy from scalar and vector zero point fields may not apply to *primary interactions*; but the *secondary interactions* of QED, QCD etc do apply in the usual manner to complete infinite superpositions.

Finally, if the fundamental particles can in fact be built from infinite superpositions this suggests:

- *Quantum mechanics rules the expansion of space and the warping of spacetime around concentrations of mass/energy, but only if it is in the form of infinite superpositions that cutoff* $@\lambda\!\!\!^{-}_{\max} \approx R_{CCH}$.
- *The expansion of space and warping of spacetime may possibly be the only direct evidence of infinite superpositions we will ever see.*
- *General Relativity is a consequence of Quantum Mechanics; it is a local effect, but to vast scales. Like all long range forces it does not reach past the horizon radius* R_{CCH} .
- *General Relativity may not apply inside Black Holes, and the event horizon itself may very well be a Spacetime Boundary.*

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