# Spontaneous free-particle acceleration in quantum electrodynamics with a real electromagnetic zero-point field

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It is shown that if the electromagnetic zero-point field (ZPF) of quantum electrodynamics (QED) is taken in the realistic sense, at least within the nonrelativistic approximation of QED, free unconfined electromagnetically interacting particles increase their translational kinetic energies when exclusively submitted to the action of the ZPF. This prediction is assessed and compared with other well-known difficulties associated with a realistic version of the ZPF concept. A parallel with an analogous acceleration phenomenon predicted in stochastic electrodynamics (SED) is given and possible astrophysical applications are suggested. It is shown that neither QED, with ZPF, nor SED lead to the Schrödinger equation (SE) of ordinary quantum mechanics, despite interesting arguments to the contrary, due precisely to this acceleration prediction.

## I. INTRODUCTION

The zero-point field (ZPF) results from quantizing the electromagnetic cavity modes in the "höhlraum." It leads to experimentally confirmed predictions like the Casimir effect[1] and several others well known in quantum optics.[2] Despite these results, when the ZPF is taken in a realistic sense, unsurmounted gravitational and thermodynamical difficulties[3] ensue that historically[4] have led to the denial of any reality for the ZPF which since then has been taken merely as a virtual field.[5] Here we study an additional difficulty that results from taking the ZPF as a real background field within quantum electrodynamics (QED) (at least in its nonrelativistic approximation). This effect is that free electromagnetically interacting particles (monopolar or polarizable) when exclusively left to the action of the ZPF background are seen to spontaneously increase their translational kinetic energy. This acceleration phenomenon is known to occur in classical electrodynamics with a classical electromagnetic ZPF[6–10] also called random, or more commonly, stochastic electrodynamics (SED).[11] Intuitively we can see this happening in a classical sense because due to the Lorentz invariance of the energy density spectrum of the ZPF[6,12,13] there is no possible Einstein-Hopf drag force on a particle moving through the ZPF. This cancellation seems to occur for all possible internal configurations of the particle.[14] On the other hand, the combined effect of the electric and magnetic fields of the ZPF background on the particle produce a fluctuating impulsive effect that in its turn generates a now unchecked energy growth. The impulse results from the action of forces somewhat reminiscent of those producing the radiation pressure on the electrons of a metallic surface. The energy growth is unchecked only for the ZPF background due in this case to the cancellation of the compensating Einstein-Hopf dissipation mentioned above.

Only two effects have so far been considered that partially quench the ZPF acceleration effect of SED. The first is collisions (with the walls[6,15] of a containing cavity and interparticle collisions). The second[16] is more fundamental. It refers to the internal structure of the accelerated particle: if the center of charge and the center of mass of the particle are not coincident but admit relative displacements and if the coupling between the two centers allows for a ZPF-induced relativistic motion of the center of charge around the center of mass, we have a model that amounts to a classical version of the Zitterbewegung[17] of ordinary QED. This Zitterbewegung quenches partially or totally the acceleration.[16] This quenching effect may easily be extended to other than simple monopolar particles, e.g., the partial quenching in polarizable particles.[18] In this paper we consider both monopolar and polarizable particles, but as the treatment is nonrelativistic we omit any Zitterbewegung considerations and hence any possible resulting partial quenching of the acceleration effect.

A final point to make is that the existence of the acceleration effect, both in QED (at least in nonrelativistic QED) with a ZPF and in SED, establishes a connection between these theories as well as a separation of both of them from ordinary quantum mechanics (QM). In particular, we claim that the acceleration effect is foreign to the behavior predicted by the Schrödinger equation (SE) for a free particle, because the SE predicts that the free particle remains in its original energy eigenstate. There have been several interesting attempts at deriving the SE from ordinary SED.[19–21] Inspection of those attempts, however, shows that the authors have in a nonrelativistic approximation neglected the effect of the magnetic field in the Lorentz force. Nevertheless, as pointed out above, it is precisely the cross effect of the electric and the magnetic fields of the ZPF in the expression for the Lorentz force on the particle that gives the forces that produce the acceleration effect.

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## II. TRANSLATIONAL ENERGY GROWTH

The particle acceleration of SED has been[7–11,14–16] studied by means of the Abraham-Lorentz equation. One considers a frame where the particle is moving slowly or where it is instantaneously at rest. The extension to relativistic speeds comes easily because the rate of translational energy growth $\Omega \equiv dE/dt$ is an invariant. Thus, in order to check if there is also an acceleration under the ZPF of QED it is easier than attempting the standard perturbative approach to look for the natural Abraham-Lorentz extension to QED of the classical Abraham-Lorentz equation[22] and proceed from there by means of analogous techniques to those used in SED to prove the acceleration effect. A Heisenberg-representation operator equation[23] that corresponds to the generalized classical Abraham-Lorentz equation has recently been proposed,[22] which presupposes a nonrelativistic charged particle of mechanical mass $m_0$ with spherically symmetric charge distribution. If there is a subset of states in the Hilbert space of the problem such that the matrix elements of the square of the velocity operator $\langle m \,|\, \dot{\vec{R}}^2/c^2 \,|\, n\rangle$ are negligible, we can obtain the operator equation[22]

$$m_0\ddot{\vec{R}}(t)=e\vec{E}+\frac{e}{2c}[\dot{\vec{R}}\times\vec{B}-\vec{B}\times\dot{\vec{R}}] -\frac{2e^2}{3c^2}\sum_{n=0}^{\infty}\frac{(-1)^n}{n!c^n}A_n\frac{d^{n+2}\vec{R}(t)}{dt^{n+2}}\,, \tag{1}$$

where $\vec{R}$ is the position operator, $\dot{\vec{R}}\equiv d\vec{R}/dt$, and $\ddot{\vec{R}}=d^2\vec{R}/dt^2$. The $\vec{E}$ and $\vec{B}$ field operators represent only the incoming fields (not the self-fields). The $A_n$ are rather complicated functions of the charged-particle Compton wavelength $\lambda\!\!\!^{-}_c=\hbar/m_0c$, which we do not need to write in full generality. Equation (1) is formally very similar to the generalized Abraham-Lorentz equation of classical electrodynamics. This approach is called *nonrelativistic* QED.[22] As the electrostatic self-energy of a point charge is found to be zero[22] ($A_0=0$), we can identify the mechanical mass $m_0$ with the observed mass $m$, $m_0=m$ in Eq. (1) above. In a first approximation we neglect the effect of the magnetic field[9] and have

$$\ddot{\vec{R}}(t)+c\Gamma\sum_{n=1}^{\infty}\frac{(-1)^n}{n!c^n}A_n\frac{d^{n+2}\vec{R}(t)}{dt^{n+2}}\cong\frac{e}{m}\vec{E}\,, \tag{2}$$

where $\Gamma\equiv 2e^2/3mc^3$ is the Abraham-Lorentz time parameter and where the $\vec{E}$ field in the dipole approximation can be written in the form $\vec{E}=\vec{E}^{+}+\vec{E}^{-}$, with $\vec{E}^{-}=[\vec{E}^{+}]^{\dagger}$. The symbol $\dagger$ signifies Hermitian conjugation and

$$\vec{E}^{+}=i\sum_{l,\sigma,\eta}\left[\frac{\hbar\omega_l}{2V}\right]^{1/2}\hat{\epsilon}_{l\sigma}a_{l\sigma\eta}e^{-i\eta\omega_l t} \tag{3}$$

with the commutation relations on the $a_{l\sigma\eta}$'s given in Ref. 24. The $\eta$, $\eta=\pm1$, is introduced in the phases for convenience because after the dipole approximation we like to distinguish between incoming and outgoing waves for a given point ($\vec{k}_l$ and $-\vec{k}_l$). Accordingly, the particle's position operator $\vec{R}$ is written as $\vec{R}=\vec{R}^{-}+\vec{R}^{+}$, $R^{-}\equiv[R^{+}]^{\dagger}$, where

$$\vec{R}^{+}=i\sum_{l,\sigma,\eta}\vec{R}^{+}_{l\sigma\eta}(\omega_l)e^{-i\eta\omega_l t}\,. \tag{4}$$

Replacing in (2) we obtain

$$\vec{R}^{+}_{l\sigma\eta}(\omega_l)=-i\frac{f(\eta\omega_l)}{\eta\omega_l}\frac{e}{m}\left[\frac{\hbar\omega_l}{2V}\right]^{1/2}\hat{\epsilon}_{l\sigma}a_{l\sigma\eta}\,, \tag{5}$$

where

$$f(\eta\omega_l)=\left\{i\eta\omega_l\left[1+c\Gamma\sum_{n=1}^{\infty}\left[\frac{i\eta\omega_l}{c}\right]^n\frac{A_n}{n!}\right]\right\}^{-1} \tag{6}$$

and thus expressions for $\vec{R}^{+}$ and $\dot{\vec{R}}^{+}$ and their conjugates follow immediately. We know then the particle velocity $\dot{\vec{R}}=\dot{\vec{R}}^{+}+\dot{\vec{R}}^{-}$. The magnetic field $\vec{B}$, within the dipole approximation, is $\vec{B}=\vec{B}^{+}+\vec{B}^{-}$, $\vec{B}^{-}=[\vec{B}^{+}]^{\dagger}$ and

$$\vec{B}^{+}=i\sum_{l,\sigma,\eta}\left[\frac{\hbar\omega_l}{2V}\right]^{1/2}(\hat{k}_l\times\hat{\epsilon}_{l\sigma})a_{l\sigma\eta}e^{-i\eta\omega_l t}\,. \tag{7}$$

We proceed to estimate the force due to the magnetic field $(e/2c)[\dot{\vec{R}}\times\vec{B}-\vec{B}\times\dot{\vec{R}}]$. It is a factor $\dot{\vec{R}}/c$ "smaller" than the force due to the electric field, i.e., in the subset of the Hilbert space we are interested in, its contribution is much smaller than that of the electric field: $|\langle n\,|\,e\vec{E}\,|\,m\rangle| \gg |\langle n\,|\,(e/2c)[\dot{\vec{R}}\times\vec{B}-\vec{B}\times\dot{\vec{R}}]\,|\,m\rangle|$. One obtains then

$$\begin{aligned}\frac{e}{2c}[\dot{\vec{R}}\times\vec{B}-\vec{B}\times\dot{\vec{R}}]=\frac{e^2}{2mc}\sum_{(1)}\sum_{(2)}\left[\frac{\hbar\omega_1}{2V}\right]^{1/2}\left[\frac{\hbar\omega_2}{2V}\right]^{1/2}\\ \times\hat{\epsilon}_1\times(\hat{k}_2\times\hat{\epsilon}_2)[f_1e^{-i(\eta_1\omega_1+\eta_2\omega_2)t}(a_1a_2+a_2a_1)-f_1e^{-i(\eta_1\omega_1-\eta_2\omega_2)t}(a_1a_2^{+}+a_2^{+}a_1)\\ -f_1^{*}e^{-i(-\eta_1\omega_1+\eta_2\omega_2)t}(a_1^{+}a_2+a_2a_1^{+})\\ +f_1^{*}e^{-i(-\eta_1\omega_1-\eta_2\omega_2)t}(a_1^{+}a_2^{+}+a_2^{+}a_1^{+})]\,,\end{aligned} \tag{8}$$

where the subscript $i$, $i=1,2$, under the summation sign replaces $l_i$, $\sigma_i$, $\eta_i$ and the subscript $i$ under the letters replaces the subindices $l_i$, $\sigma_i$, $\eta_i$ (or just $l_i$ in the case of the wave vectors $\vec{k}_{l_i}$ and the frequencies $\omega_{l_i}$). Next we estimate the impulses $\vec{\Delta}=\vec{\Delta}_E+\vec{\Delta}_B$, where $\vec{\Delta}_E\equiv\int_{-\tau/2}^{\tau/2}e\vec{E}\,dt$ and $\vec{\Delta}_B\equiv(e/2c)\int_{-\tau/2}^{\tau/2}[\dot{\vec{R}}\times\vec{B}-\vec{B}\times\dot{\vec{R}}]dt$. Next we should square and average over the vacuum state $|0\rangle\equiv|n_{l\sigma\eta}=0,\forall l,\sigma,\eta\rangle$. It is a simple matter to show that after quantum mechanical averaging, $\vec{\Delta}_E^2$ gives a constant contribution very similar to that of the jiggling motion of SED[9,16] and that we identified with the transverse self-energy of the electron under the electromagnetic fluctuations of the vacuum.[9,16] The average $\langle 0|\vec{\Delta}_E\cdot\vec{\Delta}_B|0\rangle$ vanishes. We proceed then to estimate

$$\vec{\Delta}_B=\frac{e}{2c}\int_{-\tau/2}^{\tau/2}[\dot{\vec{R}}\times\vec{B}-\vec{B}\times\dot{\vec{R}}]dt$$
$$=\tau\frac{e^2}{2mc}\sum_{(1)}\sum_{(2)}\left[\frac{\hbar\omega_1}{2V}\right]^{1/2}\left[\frac{\hbar\omega_2}{2V}\right]^{1/2}$$
$$\times\hat{\epsilon}_1\times(\hat{k}_2\times\hat{\epsilon}_2)\left\{[(a_1a_2+a_2a_1)f_1+(a_1^\dagger a_2^\dagger+a_2^\dagger a_1^\dagger)f_1^*]\frac{\sin\left[\dfrac{\eta_1\omega_1+\eta_2\omega_2}{2}\tau\right]}{\left[\dfrac{\eta_1\omega_1+\eta_2\omega_2}{2}\tau\right]}\right.$$
$$\left.-[(a_1a_2^\dagger+a_2^\dagger a_1)f_1+(a_1^\dagger a_2+a_2a_1^\dagger)f_1^*]\frac{\sin\left[\dfrac{\eta_1\omega_1-\eta_2\omega_2}{2}\tau\right]}{\left[\dfrac{\eta_1\omega_1-\eta_2\omega_2}{2}\tau\right]}\right\}. \tag{9}$$

Next we estimate the rate of energy growth $\Omega\equiv dE/dt$. We have to square the $\vec{\Delta}$ operator and find the expectation value with the vacuum field. The rate of energy growth is $(2m\tau)^{-1}\langle 0|\vec{\Delta}\cdot\vec{\Delta}|0\rangle$ which when the jiggling motion due to $\vec{\Delta}_E$ is removed, boils down to

$$\Omega=\frac{1}{2m\tau}\langle 0|\vec{\Delta}_B\cdot\vec{\Delta}_B|0\rangle$$
$$=\left[\frac{1}{2m\tau}\right]\left[\frac{e^2}{2mc}\right]\tau^2\sum_{(1)(2)(3)(4)}\left[\frac{\hbar\omega_1}{2V}\right]^{1/2}\left[\frac{\hbar\omega_2}{2V}\right]^{1/2}\left[\frac{\hbar\omega_3}{2V}\right]^{1/2}\left[\frac{\hbar\omega_4}{2V}\right]^{1/2}\hat{\epsilon}_1\times(\hat{k}_2\times\hat{\epsilon}_2)\cdot\hat{\epsilon}_3$$
$$\times(\hat{k}_4\times\hat{\epsilon}_4)\left\{4(\delta_{13}\delta_{24}+\delta_{14}\delta_{23})\frac{\sin\left[\dfrac{\eta_1\omega_1+\eta_2\omega_2}{2}\tau\right]}{\left[\dfrac{\eta_1\omega_1+\eta_2\omega_2}{2}\tau\right]}\frac{\sin\left[\dfrac{\eta_3\omega_3+\eta_4\omega_4}{2}\tau\right]}{\left[\dfrac{\eta_3\omega_3+\eta_4\omega_4}{2}\tau\right]}\right\}, \tag{10}$$

where we used the fact that

$$\langle 0|(a_1a_2+a_2a_1)(a_3^\dagger a_4^\dagger+a_4^\dagger a_3^\dagger)|0\rangle=4(\delta_{13}\delta_{24}+\delta_{14}\delta_{23})$$

and observe that all other expectations give identically zero or involve the factor $\delta_{12}\delta_{34}$ which when multiplied by $\hat{\epsilon}_1\times(\hat{k}_2\times\hat{\epsilon}_2)\cdot\hat{\epsilon}_3\times(\hat{k}_4\times\hat{\epsilon}_4)$ yield $\hat{\epsilon}_1\times(\hat{k}_1\times\hat{\epsilon}_1)\cdot\hat{\epsilon}_3\times(\hat{k}_3\times\hat{\epsilon}_3)=\hat{k}_1\cdot\hat{k}_3=\hat{k}_{l_1\sigma_1}\cdot\hat{k}_{l_3\sigma_3}$ that cancels after adding over the indices because of symmetry. This can also be seen by passing to the continuum:

$$\sum_{l,\sigma}\{\cdots\}=\frac{V}{(2\pi)^3}\sum_{\sigma=1}^{2}\int d^3k\{\cdots\}$$
$$=\frac{V}{(2\pi c)^3}\sum_{\sigma=1}^{2}\int\omega^2d\omega\int d\Omega_{\hat{k}}\{\cdots\}.$$

From this we also obtain the angle integrations

$$\int d\Omega_1\int d\Omega_2\sum_{\sigma_1=1}^{2}\sum_{\sigma_2=1}^{2}\{\hat{\epsilon}(\hat{k}_1,\sigma_1)\times[\hat{k}_2\times\hat{\epsilon}(\hat{k}_2,\sigma_2)]\}^2=\tfrac{8}{3}(4\pi)^2, \tag{11}$$

where the change in notation is self-explanatory, and

$$\int d\Omega_1\int d\Omega_2\sum_{\sigma_1=1}^{2}\sum_{\sigma_2=1}^{2}\hat{\epsilon}(\hat{k}_1,\sigma_1)\times[\hat{k}_2\times\hat{\epsilon}(\hat{k}_2,\sigma_2)]\cdot\hat{\epsilon}(\hat{k}_2,\sigma_2)\times[\hat{k}_1\times\hat{\epsilon}(\hat{k}_1,\sigma_1)]=0 \tag{12}$$

because of symmetry in the integrations. We obtain then

$$\Omega=\frac{1}{(2m\tau)}\left[\frac{e^2}{2mc}\right]^2\tau^2\left[\frac{\hbar}{2V}\right]^2\left[\frac{V}{(2\pi c)^3}\right]^2\tfrac{8}{3}(4\pi)^2\int\omega_1^2d\omega_1\int\omega_2^2d\omega\sum_{\eta_1\eta_2}4\,|f_1|^2\left[\frac{\sin\left[\frac{\eta_1\omega_1+\eta_2\omega_2}{2}\tau\right]}{\left[\frac{\eta_1\omega_1+\eta_2\omega_2}{2}\tau\right]}\right]^2 . \quad (13)$$

Now we use the identity $\int_{-\infty}^{\infty}(\sin\beta/\beta)^2 d\beta=\pi$ whose argument around $\beta=0$ behaves approximately as a $\delta$ function. According to this only the terms in (13) where $\eta_1=-\eta_2$ give a substantial contribution. We obtain then

$$\Omega=\frac{\alpha}{2\pi^3}\int_0^{\infty}\left[\frac{\hbar\omega}{mc^2}\right]^2(\Gamma_m\omega)(\hbar\omega)\,|\,\omega f(\omega)\,|^2 d\omega\ , \quad (14)$$

where $\alpha=e^2/\hbar c$ and $\Gamma_m=2e^2/3mc^3$. So a monopolar particle is seen to increase its translational kinetic energy. It remains to be seen if a fully relativistic calculation still gives the energizing property, in particular, because of our discussion on Zitterbewegung, as there are reasons to think that the acceleration may be quenched for electrons. However, this is not the case for polarizable particles, for which the rate of kinetic-energy growth is only decreased. It is interesting then to give the expression for $\Omega$ in the case of a polarizable particle moving through the ZPF of QED. The analysis is very similar to the one that leads to (14). We consider a polarizable particle composed of two subparticles of equal and opposite charges $\pm e$ and of masses $M$ and $m$, respectively, but with $M \gg m$. A harmonic binding is assumed. The effect of the magnetic field of the ZPF on the big mass may be neglected and if we are not interested in the electric-field-induced jiggling motion on the particle as a whole, we may also neglect the action of the $E$ field on the large mass. We are left with

$$M\ddot{\vec{R}}+m\omega_c^2\vec{R}=m\omega_c^2\vec{r}\ , \quad (15)$$

$$m\ddot{\vec{r}}+m\omega_c^2\vec{r}=m\omega_c^2\vec{R}+e\vec{E}+\frac{e}{2c}(\vec{r}\times\vec{B}-\vec{B}\times\vec{r}) -\frac{2e^2}{3c^2}\sum_{n=0}^{\infty}\frac{(-1)^n}{n!c^n}A_n\frac{d^{n+2}\vec{r}(t)}{dt^{n+2}}\ , \quad (16)$$

where $\vec{R}$ and $\vec{r}$ are the position operators for the large and the small masses, respectively. The coupled operator equations can easily be reduced to a single equation for the $\vec{r}$ operator and an analogous procedure to that presented for the monopolar particle can be followed. We obtain again an $\Omega$ given by (14) but where $\Gamma_m$ instead of being written with $m$ is written with $M$, $\Gamma_M=2e^2/3Mc^3$ and $f(\omega)$ is given by

$$f(\omega)\equiv\left\{i\omega\left[1-\delta+c\Gamma_m\sum_{n=1}^{\infty}\left[\frac{i\omega}{c}\right]^n\frac{A_n}{n!}\right]\right\}^{-1}, \quad (17)$$

where $\delta=\omega_c^2(\omega^2-m\omega_c^2/M)^{-1}$ and $\Gamma_m=2e^2/3mc^3$ is the Abraham-Lorentz time constant for the subparticle of mass $m$. One interesting feature of Eq. (14) and its counterpart for a polarizable particle is that we obtain a convergent expression even in the point-particle limit. Moniz and Sharp[22] have obtained a closed-form solution for the terms of the series $A_n$ in the point-charge limit

$$A_{2s+1}=(-1)^s\frac{(2s+1)(8s+9)}{3(s+1)(2s+3)}[(4s+1)!!]\lambda_c^{2s}\ , \quad (18a)$$

$$A_{2s}=0\ , \quad (18b)$$

where $s=0,1,\ldots$ and $\lambda_c=c/\omega_c=\hbar/mc$ is the Compton wavelength associated with a particle of mass $m$. From (14), (18a), and (18b) we obtain

$$\Omega=\frac{\alpha}{2\pi^3}\int_0^{\infty}d\omega\left[\frac{\hbar\omega}{mc^2}\right]^2(\Gamma_M\omega)(\hbar\omega)\left\{(1-\delta)^2+(\Gamma_m\omega)^2\frac{1}{3}\left[\sum_{s=0}^{\infty}\left[\frac{\omega}{\omega_c}\right]^{2s}\frac{(8s+9)}{(s+1)(2s+3)}\frac{(4s+1)!!}{(2s)!}\right]^2\right\}^{-1}, \quad (19)$$

where we have written the expression for the case of a polarizable particle with $\delta$ as defined above and $\Gamma_M\neq\Gamma_m$. The case of a monopolar particle, Eq. (14), is recovered by setting $\delta=0$ and $\Gamma_M=\Gamma_m=2e^2/3mc^3$ where $m$ is the mass of the monopolar particle itself. In the classical case it was found[9] that in the point-charge limit ($\omega_m\to\infty$) $\Omega$ diverges, but this was due to the fact that we did not use the full-radiation-reaction equation[25] which, as in Eq. (1), contains a series that ultimately guarantees the convergence.

## III. DISCUSSION

From the argument in Sec. I it is clear that ordinary SED cannot lead to the Schrödinger equation and hence that ordinary SED is not a theory that necessarily implies QM. When, on the grounds of a nonrelativistic approximation, the magnetic field is dismissed in the equation of Abraham-Lorentz in SED, one does not obtain the acceleration prediction but only a jiggling motion. The ensuing nonaccelerative scheme is what permits one, after reasonable approximations, to derive a Schrödinger-type equation within SED.[11] This is reminiscent of the Fenyes-Nelson[26] derivation of the Schrödinger equation in stochastic mechanics (SM) where only a random field is considered, which produces a random motion that ultimately leads to the "nonclassical" behavior of the particle. The random field of SM is very general. It does not even need to be electromagnetic. In this sense one usually sees

SM as a more general theory than SED. In SED the random field is very special. It has the unique Lorentz-invariant spectral energy density for a random electromagnetic (EM) field and thence it lacks the usual dissipative feature of the Einstein-Hopf drag force.[14] The only dissipation left, to compensate for the induced ZPF fluctuations, is that due to the radiation reaction. In QED there also is acceleration when the QED-predicted vacuum fluctuations of the EM field are considered in the real sense and are included as a random background field in the problem of the free particle. So, QED with its ZPF not dismissed is not consistent with QM. This inconsistency is just another difficulty that adds itself up to other divergencies[27] so common in QED.

The ZPF acceleration presents some thermodynamic difficulties. A free particle is seen to accelerate spontaneously. This seems to violate the first law, but indeed it does not, as the ZPF, when taken as real, has an infinite amount of energy. The no violation of the second law is less clearly explained. A gas of infinite noncolliding charges in an infinite unbound space is seen to increase its translational kinetic energy by extracting energy from the ZPF. It is as if energy were spontaneously transferred from a reservoir at zero temperature to a reservoir at a higher temperature without the provision of any work. But one of the premises involved in this objection is not correct. If the ZPF, at least that of SED,[11] is originated in the motion of all accelerated charges in the universe, one cannot claim that it represents a background radiation at zero temperature but only that it is the radiation that is left when thermal radiation is removed. This ZPF radiation is the vehicle by which the given free particle tends to establish its equilibrium with the rest of the universe. Recall that in SED the ZPF has a fundamental property that transcends the ordinary phenomenology of random EM fields in the thermodynamics of thermal radiation.[28] There is also the following question: if free particles are spontaneously accelerated by the ZPF, why is it that there is not a spontaneous warming up of ordinary matter in the universe? A qualitative answer may be given. It contains two complementary aspects: (i) When matter fields are strong, $\rho(\omega)$ does not have the Lorentz-invariant distribution and drag forces are present; (ii) The cooling effect of collisions makes the particles radiate back energy into the random-field background. This shows that thermodynamic difficulties are not compelling objections to the reality of the ZPF concept. On the positive side we mention the acceleration mechanism proposed as a possible source of excitation of the IGS plasma and as a possible originator of the x-ray background and of cosmic rays (CR) primaries preferentially in the ultrahigh energies ranges.[7,18,19] The proposal was made considering only the SED original acceleration.[6] However, all still holds if the mechanism can be rephrased within QED. As the astrophysical aspects have been discussed at length elsewhere,[7,18,29] we omit further comments on them. But before accepting the QED acceleration as an established fact, several difficulties should first be surmounted: (i) Nonrelativistic QED is not Lorentz invariant and a Lorentz-invariant QED treatment should be provided for the acceleration prediction; (ii) The presentation of Moniz and Sharp[22] does not follow the standard perturbative recipes of QED but introduces different approximations whose predictions are sometimes at variance with those of the usual perturbative analysis[30] and a comparison between the traditional and this nonperturbative approach must first be accomplished; (iii) Equation (1) neglects the operator $\dot{\vec{R}}^2/c^2$ without providing a conclusive mathematical justification for this approximation;[22] (iv) In our derivation we use the dipole approximation and in SED this approximation is reasonable[8] but we have to better justify its use in QED; (v) In our proof of the acceleration in QED we use an iterative procedure where first we neglect the magnetic field and obtain $\dot{\vec{R}}$ as produced only by the action of the electric field. Next we introduce this $\vec{R}$ into the magnetic part of the Lorentz force operator and proceed to obtain the impulse operator $\vec{\Delta}_B$ and then we obtain the acceleration. This is the standard procedure in SED but its implementation in QED requires more rigor.

Finally we observe that a classical form of Zitterbewegung is known to curb down or even quench the classical acceleration predicted in SED.[16] This suggests that electron Zitterbewegung in QED may quench the acceleration predicted here for structureless monopolar charges in the nonrelativistic QED of Moniz and Sharp. A fully relativistic QED calculation which de facto would imply Zitterbewegung might not display the acceleration effect as suggested by the quenching of the acceleration for a classical but relativistic model of a monopolar charged particle where the center of charge and the center of mass do not coincide,[16] but for polarizable particles there is still acceleration.[18]


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