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Ekman, R

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Reduction of order, resummation and radiation reaction

Robin Ekman,^{1,*} Tom Heinzl,^{1,†} and Anton Ilderton^{1,‡}

¹*Centre for Mathematical Sciences, University of Plymouth, Plymouth, PL4 8AA, UK*

The Landau-Lifshitz equation is the first in an infinite series of approximations to the Lorentz-Abraham-Dirac equation obtained from ‘reduction of order’. We show that this series is divergent, predicting wildly different dynamics at successive perturbative orders. Iterating reduction of order *ad infinitum* in a constant crossed field, we obtain an equation of motion which is free of the erratic behaviour of perturbation theory. We show that Borel-Padé resummation of the divergent series accurately reproduces the dynamics of this equation, using as little as two perturbative coefficients. Comparing with the Lorentz-Abraham-Dirac equation, our results show that for large times the optimal order of truncation typically amounts to using the Landau-Lifshitz equation, but that this fails to capture the resummed dynamics over short times.

I. INTRODUCTION

Radiation reaction (RR) in electrodynamics becomes relevant in the presence of strong fields, where RR forces can become comparable to, or dominate, the Lorentz force. In the presence of strong gravitational fields RR was observed decades ago in studies of the Hulse-Taylor binary pulsar [1]. The recent observation of gravitational waves and their analysis has further highlighted the importance of RR e.g. at third post-Minkowskian order [2], triggering a number of investigations to clarify the subtleties involved [3–5]. The direct connection between RR in gravity and in Yang-Mills theory has also recently been explored in the context of double copy [6–8]. Here we consider RR in classical electrodynamics [9–11], where it was first formulated theoretically. Our interest is motivated by emerging experimental access to previously uncharted strong-field regimes, provided by intense lasers [12–18].

The classical equation of motion supposed to describe RR is the third-order Lorentz-Abraham-Dirac (LAD) equation [9–11]. Its third-order character is in conflict with Newton-Laplace determinism, as more than two initial conditions are needed to uniquely determine a solution. Adding an initial condition for acceleration makes the initial-value problem well posed, but leads to unphysical runaway solutions at temporal infinity. Imposing instead Dirac’s condition of vanishing final acceleration [11], one has an initial-boundary value problem, and solutions (the existence and uniqueness of which is not guaranteed [19–22]) exhibit pre-acceleration before the external field is encountered (albeit on a small time scale of about 2 fm/c). Finally, both analytical and numerical solutions of LAD are hampered by the strong nonlinearities present in the fully relativistic case [23].

In view of these difficulties, it is common to adopt the method of ‘reduction of order’. When applied to LAD, this yields the Landau-Lifshitz (LL) equation [24], which, being second-order in time derivatives, yields a well-posed initial-value problem and hence is free of runaway and pre-accelerating solutions [25]. Reduction of order from LAD to LL is an example of singular perturbation theory [26]. It is known that the ensuing perturbative approximations to the full equation will typically miss nonlinear phenomena such as layers and bifurcation branches unless one introduces an appropriate amount of parameter fine-tuning [19–22, 27]. Furthermore, LL represents only the first of an infinite series of approximations to LAD obtained by *iteration* of reduction of order.

With this in mind, we show here that the perturbation expansion generated by iterating reduction of order has zero radius of convergence and is asymptotic, predicting wildly different physics at successive orders of perturbation theory. We are therefore prompted to investigate resummation of the series, as has recently been highlighted in investigations of both classical [28] and quantum [29] radiation reaction, and in strong field problems more generally [30–33].

We focus mainly on the case of a constant crossed field (CCF) background, which is relevant to the Ritus-Narozhny conjecture on the breakdown of perturbation theory in strong fields [34]. We will show that the simplifications associated with the CCF case allow for an almost complete resummation of the divergent perturbative series: we are able to iterate reduction of order an *infinite* number of times, casting the result as a nonlinear system of first-order, ordinary differential equations (ODEs). While the perturbative expansion of this system is divergent, we are able to prove that the strong-field expansion is convergent, a situation reminiscent of the analogous expansions of the Heisenberg-Euler Lagrangian [35].

Furthermore, the system can be solved numerically to high accuracy, such that the solution may be viewed as numerically ‘exact’. We compare this exact result to resummed Padé-Borel approximants of the perturbative expansion,

* robin.ekman@plymouth.ac.uk

† theinzl@plymouth.ac.uk

‡ anton.ilderton@plymouth.ac.uk

showing that these can match the numerical answer to high precision using only a few terms. We find that infinite reduction of order gives results which are free of the erratic behaviour seen in perturbation theory, and which match LAD, but without the non-perturbative, and unphysical runaways or pre-acceleration behaviour. All our results indicate that radiation reaction rapidly drives particle motion to a regime where LL is valid; the optimal order of truncation for the divergent reduction of order expansion is therefore one.

We work in the classical theory throughout; quantum corrections are typically expected [36–38] when high field strengths cause hard acceleration gradients and significant radiation reaction, but the inclusion of such corrections goes beyond the scope of this paper. Our focus will therefore be on resummation and physical properties of radiation reaction equations, rather than on phenomenological predictions.

This paper is organised as follows. In Sec. II we introduce reduction of order, and its iterations, of the relativistic LAD equation in a constant crossed field of arbitrary strength. We show that the resulting series is divergent, highlight how this manifests in the physics of particle motion, and how this can be cured using Borel-Padé resummation. In Sec. III we iterate reduction of order to all orders, obtaining an equation of motion that we dub LL_∞ . We show that the strong-field expansion of LL_∞ is convergent and use this to investigate the physics of the strong field regime. In Sec. IV we use the convergent strong-field behaviour to improve our perturbative resummation, and obtain an analytical expression for LL_∞ . We then compare this with the numerical solution of LAD. We conclude in Sec. VI, suggesting a physical explanation of the divergence of perturbation theory, relating our results to the previous literature, and discussing extensions. Throughout we use units where $c = \epsilon_0 = \hbar = 1$.

II. ITERATION OF REDUCTION OF ORDER

Consider the orbit $x^\mu(\tau)$ of a particle, charge e and mass m , in an external field $F_{\mu\nu}(x)$. The orbit is parameterised by proper time τ , the particle velocity is written $u^\mu = \dot{x}^\mu$, and overdots denote derivatives with respect to τ . The LAD equation of motion is, writing $f_{\mu\nu} = eF_{\mu\nu}/m$,

$$\dot{u}^\mu = f^{\mu\nu}u_\nu + \tau_0 \mathcal{P}^{\mu\nu}\ddot{u}_\nu, \quad (1)$$

in which $\mathcal{P}^{\mu\nu} = \eta^{\mu\nu} - u^\mu u^\nu$ projects orthogonally to u^μ and $\tau_0 = e^2/6\pi m = 2\alpha/3m$, α being the fine-structure constant, is the characteristic time scale of RR; for an electron $\tau_0 \approx 6.3 \times 10^{-24}$ s. Note that (1) is a third-order ODE for the orbit x^μ ; the associated difficulties of either runaways or pre-acceleration motivate the adoption of reduction of order in derivatives. This refers to formally differentiating the LAD equation, then using the equation itself to eliminate \ddot{u}^μ (and higher derivatives) in favour of new τ_0 -dependent terms. This results, in principle, in an infinite series of terms which can be truncated at any chosen order in τ_0 due to the smallness of that parameter compared to relevant timescales. Truncating at order τ_0 yields the Landau-Lifshitz equation [24, 39] (LL_1 from here, for reasons which will become clear),

$$\dot{u}^\mu = f^{\mu\nu}u_\nu + \tau_0 f^{\mu\nu,\rho}u_\nu u_\rho + \tau_0 \mathcal{P}^{\mu\nu}f_{\nu\rho}^2 u^\rho + \mathcal{O}(\tau_0^2). \quad (2)$$

Unlike (1), this is now a *second* order ODE for x^μ , hence the term ‘reduction of order’. This process can be iterated, truncating at *higher orders* in τ_0 , but rapidly becomes complicated due to the appearance of new tensor structures and higher derivatives of the field tensor.

As we are interested in large orders in τ_0 we simplify matters by limiting the number of possible terms which can appear. To this end we begin by restricting our considerations to constant backgrounds for which all field derivatives are identically zero, and we need make no approximation on them. Note that, for general fields, this may be viewed as the leading order in a derivative expansion, and there are good arguments for dropping derivative terms: these scale with powers of $\omega\tau_0$ where ω is a typical field frequency scale. For realistic fields $\omega\tau_0 \ll 1$, and indeed it can be seen in e.g. the exact solution of LL_1 in a plane wave [40–42] that the derivative terms only ever yield sub-leading effects in the field strength. Iterating reduction of order to $\mathcal{O}(\tau_0^2)$ [43, 44], it can again be shown explicitly that derivative terms have negligible impact on the physics [44]. This is also consistent with the standard effective field theory ordering of field ‘operators’ according to their dimension. This argument comes with the caveat that the derivative terms can contribute to the *non-perturbative* features of LAD [11, 23, 45], which are also of interest.

The tensor structures appearing are further simplified by taking the background to be a CCF (the zero frequency limit of a plane wave), as discussed in the introduction. The field strength is $f_{\mu\nu} = m\mathcal{E}(n_\mu\epsilon_\nu - \epsilon_\mu n_\nu)$ with n_μ lightlike, ϵ_μ spacelike, $\epsilon \cdot n = 0$, and dimensionless amplitude \mathcal{E} . Consider now iterating reduction of order with this background. Because $f_{\mu\nu}^3 \equiv 0$, there are only two tensor structures, proportional to f or f^2 , which can ever appear. For a CCF, all field invariants vanish, and the only nontrivial, dimensionless, invariant which can be constructed from the particle velocity u_μ and the field is

$$\chi = \sqrt{\frac{u^\mu f_{\mu\nu}^2 u^\nu}{m^2}} = \mathcal{E} n \cdot u, \quad (3)$$

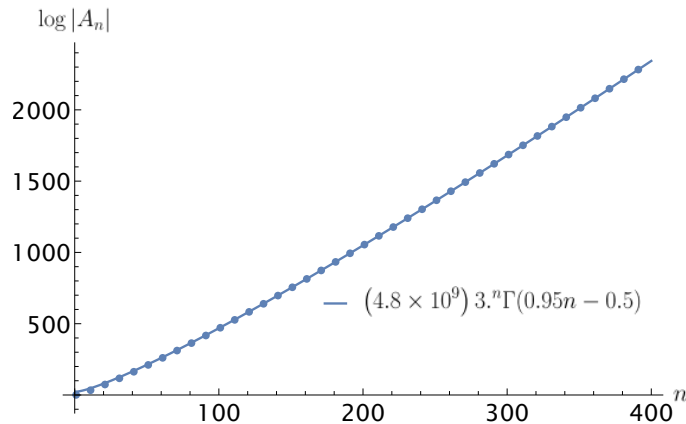


FIG. 1. The sequence A_n (circles) and a fit using the Γ function (solid). The sequence B_n has the same behaviour.

which can be interpreted as the field magnitude ‘seen’ by the particle in its instantaneous rest frame, in units of the Sauter-Schwinger field m^2/e [46, 47]. Note that χ is a composite parameter, essentially the product of field strength and particle energy.

It follows that reduction of order, iterated to *arbitrary* order in τ_0 , will yield an equation of the form

$$\dot{u}^\mu = \mathcal{A}(\chi) f^{\mu\nu} u_\nu + \tau_0 \mathcal{B}(\chi) (\mathcal{P} f^2)^{\mu\nu} u_\nu, \quad (4)$$

with the functions \mathcal{A} and \mathcal{B} depending on χ . It is easily seen that LL_1 corresponds to $\mathcal{A} = \mathcal{B} = 1$ which sets the initial condition for the iteration procedure. The equation of motion (4) is our first result. Iterated reduction of order in more general fields may be explored by applying resummation to a perturbative expansion of the integro-differential formulation of LAD, see Ref. [36, 48]

We now turn to the explicit construction of the functions \mathcal{A} and \mathcal{B} . From here on we set the electron mass m to 1 in the text, in order to simplify our equations. We re-instate m in some figures so that the reader can easily see the physical scales.

A. Perturbation theory and divergence

Iterating reduction of order, i.e. retaining terms of order up to and including τ_0^k , yields a sequence of equations of form (4) which we refer to as LL_k , $k \geq 1$. The functions \mathcal{A} and \mathcal{B} then have series expansions in $\tau_0^2 \chi^2$, given by

$$\mathcal{A} \rightarrow \mathcal{A}^{(k)} \equiv \sum_{\ell=0}^{\lfloor k/2 \rfloor} A_\ell (\tau_0 \chi)^{2\ell}, \quad \mathcal{B} \rightarrow \mathcal{B}^{(k)} \equiv \sum_{\ell=0}^{\lfloor \frac{k-1}{2} \rfloor} B_\ell (\tau_0 \chi)^{2\ell}, \quad (5)$$

in which $A_0 = B_0 = 1$, recovering the Lorentz force and LL_1 equations, respectively. After some algebra, reduction of order implies the following recursion relations for the coefficients A_k and B_k :

$$A_{k+1} = -2 \sum_{\ell=0}^k (\ell+1) A_\ell B_{k-\ell}, \quad B_{k+1} = \sum_{\ell=0}^k A_\ell A_{k-\ell} - 2 \sum_{\ell=0}^{k-1} (\ell+1) B_\ell B_{k-\ell-1}. \quad (6)$$

Calculating the first few coefficients we find

$$\begin{aligned} A_k &= \{1, -2, 20, -328, 7024, -179264, \dots\}, \\ B_k &= \{1, -6, 80, -1520, 35760, -976704, \dots\}, \end{aligned} \quad (7)$$

which are seen to alternate in sign and grow quickly. There are no matching entries in the Online Encyclopedia of Integer Sequences [49], and we are not aware of a combinatorial interpretation. Because of the factors ℓ appearing inside the sums, the coefficients can be expected to grow factorially; a simple fit confirms graphically that both sequences grow asymptotically as $3^n \Gamma(0.95n - 0.5)$, see¹ Fig. 1. The series thus has zero radius of convergence.

¹ While it is known that there is information encoded in such asymptotic growth rates [50], it is out of scope for our present purposes, as we will obtain accurate resummations with only a handful terms.

To highlight the physical implications of the divergence of perturbation theory, we present some explicit solutions to the equations of motion LL_k for low $k > 1$. To do so we use lightfront coordinates $x^\pm := x^0 \pm x^3$, $x^\perp = \{x^1, x^2\}$, choosing $n \cdot x = x^+$, lightfront time, and $\chi = \mathcal{E}u^+$. The LL_k equations of motion for u^+ and u^\perp decouple for all k , and become, for any approximation to the functions \mathcal{A} and \mathcal{B} ,

$$\frac{du^+}{dx^+} = -\tau_0 \chi^2 \mathcal{B}(\chi), \quad u^+ \frac{d}{dx^+} \frac{u^\perp}{u^+} = -\mathcal{A}(\chi) \epsilon^\perp, \quad (8)$$

with u^\perp determined by the mass-shell condition. The u^\perp components are determined by quadrature once u^+ is known, so we will focus on determining u^+ in what follows. The equation in (8) for u^+ is separable, with solution

$$\tau_0 \int_0^{x^+} dy^+ = \tau_0 x^+ = - \int_{u_0^+}^{u^+} \frac{dv}{\mathcal{E}^2 v^2 \mathcal{B}(\mathcal{E}v)}. \quad (9)$$

This *causal* integral can be performed as (for \mathcal{B} calculated in perturbation theory) the integrand is a rational function. At first and third order this yields

$$\text{LL}_1 : \quad u^+ = \frac{u_0^+}{1 + \tau_0 \mathcal{E}^2 x^+} \quad (10)$$

$$\text{LL}_3 : \quad \frac{1}{\mathcal{E}^2 v} - \frac{\sqrt{6}\tau_0}{\mathcal{E}v} \operatorname{arctanh} \sqrt{6}\mathcal{E}\tau_0 v \Big|_{v=u_0^+}^{v=u^+} = \tau_0 x^+, \quad (11)$$

with LL_5 admitting a similar, but unwieldy and unenlightening expression, which we omit. Note that, even at only third order, we just obtain an implicit expression for u^+ . We therefore proceed graphically, showing solutions to LL_1 , LL_3 , and LL_5 in Fig. 2. The features of these solutions can be read off from the respective series expansion of \mathcal{B} . Recall first that u^+ is conserved according to the Lorentz force equation, but not according to LL_1 . As $\mathcal{B}^{(3)}(\chi) = 1 - 6(\tau_0 \chi)^2 < 1$, LL_3 predicts less RR than LL_1 , up until a stationary solution where $\mathcal{B}^{(3)}$ crosses zero and hence χ , therefore u^+ , is again conserved. For χ above the zero-crossing LL_3 gives an RR force in the ‘wrong’ direction, hence predicting a runaway u^+ which goes to infinity. On the other hand LL_5 eventually predicts stronger RR than LL_1 , as $\mathcal{B}^{(5)} \rightarrow +\infty$ for large χ . This is quantitatively inconsistent with LAD; it has been shown that, writing u_0^+ for the initial velocity, $u_0^+ \geq u_{\text{LAD}}^+ \geq u_{\text{LL}_1}^+$ [51], i.e. LL_1 overestimates RR compared to LAD [48]. So, whenever $u_{\text{LL}_n}^+ < u_{\text{LL}_1}^+$, the latter is the better approximation.

The conclusions for LL_3 and LL_5 generalise: as the $\mathcal{B}^{(n)}$ alternate between diverging to $\pm\infty$, solutions to LL_n will be radically different from order to order. For orders without a runaway, quantitative agreement with LAD becomes worse and worse as $\mathcal{B}^{(n)}$ grows more rapidly with χ . We speculate that some of these features may be related to a bifurcation phenomenon associated with LAD in plane wave fields which has recently been discovered in the non-relativistic limit² [22]. In the absence of any fully relativistic solution to LAD in a plane wave, we try to extract physical results from the diverging series (5) by resumming them. We will thus obtain an equation of motion, LL_∞ , which is free of the problems of LL_k .

B. Resummation

We note that the series (5) with coefficients (7) are, although divergent, Borel summable. This prompts two questions. First, is there an optimal order of truncation (an optimal LL_k), as is typical of asymptotic series? Second, what insights can be gained from resumming the series? We investigate the second question here, returning to the first later on. We will use the Borel-Padé [52–54] method to resum the series A_k, B_k . While there are other resummation methods of potentially higher accuracy, such as Meijer- G resummation [55], Borel-Padé has been successfully applied to several topics in QED [29, 33, 56–58], is comparatively simple to implement, and we can easily generate many terms, should they be needed. (We do not expect any instabilities related to nonperturbative imaginary parts typical for non-alternating coefficients [59, 60].)

The method is as follows. Given $N + M$ perturbative coefficients A_ℓ for \mathcal{A} as a series in $(\tau_0 \chi)^2$, their Borel transform is $\sum A_\ell t^\ell / \ell!$. One constructs a Padé approximant of order N/M of the Borel transform,

$$P_{\mathcal{A}}[N/M](t) = \frac{\sum_{\ell=0}^N c_\ell t^\ell}{1 + \sum_{j=1}^M d_j t^j} = \sum_{\ell=0}^{M+N} \frac{A_\ell}{\ell!} t^\ell + \mathcal{O}(t^{M+N+1}), \quad (12)$$

² For simple potential steps similar observations go back to [19] and [20].

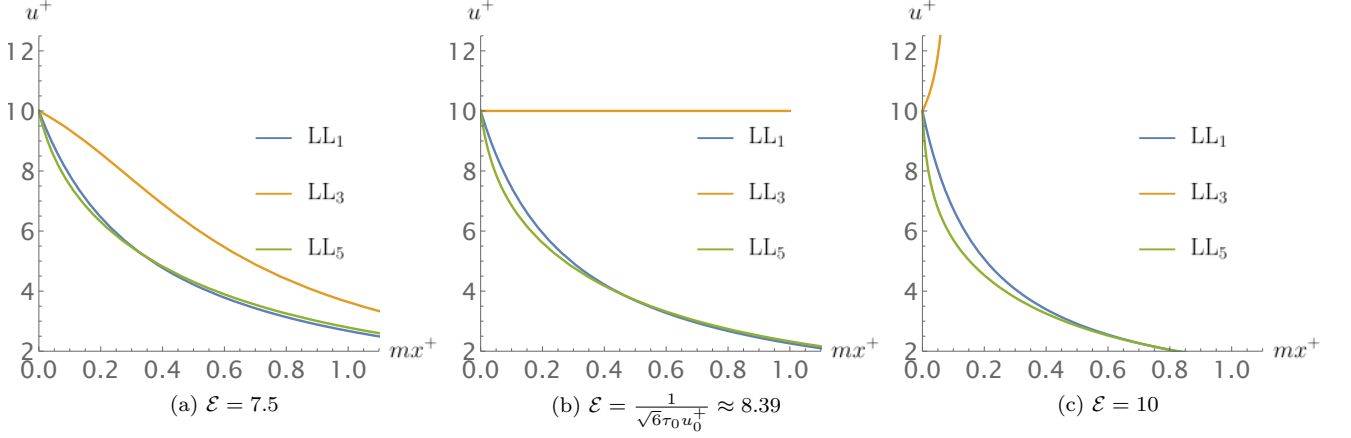


FIG. 2. Lightfront momentum according to the LL_1 , LL_3 , and LL_5 equations of motion. For a critical value of χ , LL_3 has a stationary solution; for larger χ this becomes a runaway. LL_5 is quantitatively wrong in overpredicting the strength of RR at early times, before χ becomes small.

and the resummed series is given by the inverse Borel transform

$$\mathcal{A}^{[N/M]}(\chi) := \int_0^\infty dt e^{-t} P_{\mathcal{A}}[N/M](t(\tau_0\chi)^2), \quad (13)$$

with similar expressions holding for \mathcal{B} . The rate of convergence can depend on the choice of N, M . We have found $N = M - 1$ to give the fastest convergence (and we will see in the next section why this is). However, we stress that other choices converge to the same functions, just requiring more terms to do so. The lowest-order resummants have comparatively simple analytical expressions,

$$\mathcal{A}^{[0/1]}(\chi) = \frac{e^{\frac{1}{2\tau_0^2\chi^2}} E_1\left(\frac{1}{2\tau_0^2\chi^2}\right)}{2\tau_0^2\chi^2} \xrightarrow{\chi \rightarrow \infty} \frac{2\log \tau_0\chi + \log 2 - \gamma_E}{2\tau_0^2\chi^2}, \quad (14)$$

$$\mathcal{B}^{[0/1]}(\chi) = \frac{e^{\frac{1}{6\tau_0^2\chi^2}} E_1\left(\frac{1}{6\tau_0^2\chi^2}\right)}{6\tau_0^2\chi^2} \xrightarrow{\chi \rightarrow \infty} \frac{2\log \tau_0\chi + \log 6 - \gamma_E}{6\tau_0^2\chi^2}, \quad (15)$$

where E_1 is the exponential integral [61, Ch. 5]. These give, comparing against resummants calculated using more terms, accurate results up to $\tau_0\chi \lesssim 0.3$, as shown in Fig. 3. In contrast to the perturbative series the resummants are monotonically decreasing with χ . The resummed LL_n equations of motion therefore predict less RR than LL_1 , and may have better agreement with LAD. Notably, since the resummants are always positive, the unphysical runaway solution is eliminated from those LL_n that feature it, such as LL_3 . Instead, u^+ is monotonically decreasing, which means that at large lightfront times $\chi \rightarrow 0$; as a result, $\mathcal{A}, \mathcal{B} \rightarrow 1$, so that the dynamics becomes governed by LL_1 .

The right-hand panels of Fig. 3 show that for a given χ the perturbative series agrees with the resummants when no more than $\sim \frac{1}{2}(\tau_0\chi)^{-2}$ terms are included. This is in line with the heuristic rule for asymptotic series, that the optimal truncation order is after the smallest term, which occurs at order inversely proportional to the expansion parameter [62]. As the coefficients defined by (6) grow factorially from the outset, the optimum order of truncation is one, i.e. LL_1 for $\tau_0\chi \gtrsim 0.3$. This suggests that, given their much more complicated form and \mathcal{A}, \mathcal{B} not differing greatly from *unity* for small χ , in practice no finite-order LL_k , $k \geq 1$, is ‘better’ than LL_1 . In the following we will examine the resummed, all-orders equation of motion. (We will see that the asymptotic logarithms in (14) and (15) are artefacts of the resummation procedure, but the qualitative features are as described here.)

The resummation above should be contrasted with that in the non-relativistic limit. In that case the LAD equation is linear in the field strength, resummation only involves derivative terms, and turns out to be straightforward. It can be used to recover the pre-acceleration solution of LAD from perturbative, reduction of order approximations [45].

III. LANDAU-LIFSHITZ EQUATION TO ALL ORDERS

Consider now the equation (4) obtained by iterating reduction of order infinitely many times, i.e. including (5) to all orders in τ_0 . We call this equation LL_∞ . By demanding that this be a fixed point of reduction of order, or equivalently

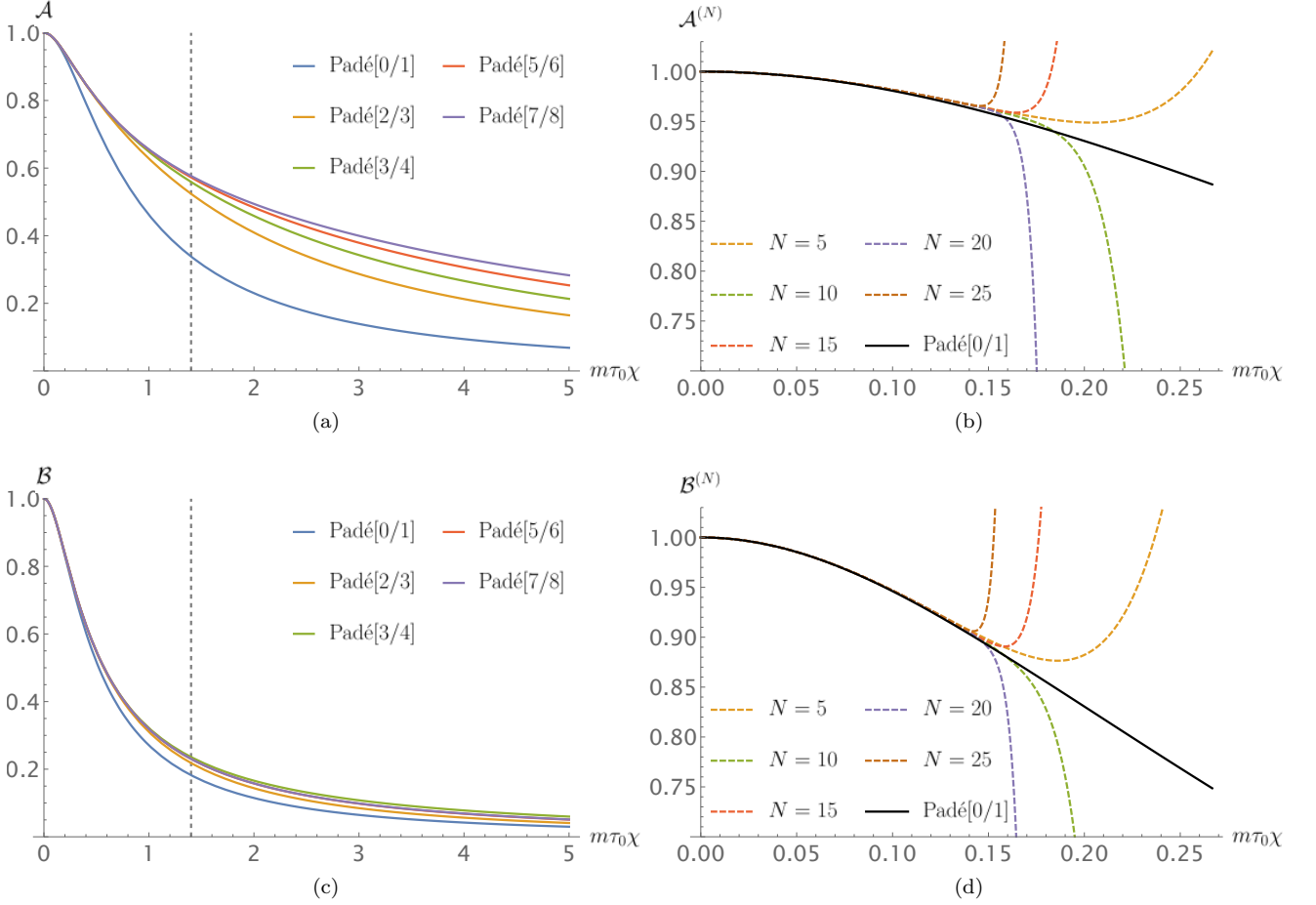


FIG. 3. Panels (a), (c): Borel-Padé resummants of the sequences A_k, B_k . The dashed vertical line marks the radius of convergence of the large- χ expansion, see Sec. III A. Panels (b), (d): Finite-order partial sums (dashed) and the lowest-order resummant (solid black). On the interval displayed in these panels one term is sufficient for convergence.

by demanding that u^μ in (4) obeys LAD, one finds that \mathcal{A} and \mathcal{B} are fully defined by the initial value problem

$$\begin{cases} \tau_0^2 \chi^3 \mathcal{B}(\chi) \mathcal{A}'(\chi) = 1 - \mathcal{A}(\chi) - 2\tau_0^2 \chi^2 \mathcal{A}(\chi) \mathcal{B}(\chi), \\ \tau_0^2 \chi^3 \mathcal{B}(\chi) \mathcal{B}'(\chi) = -\mathcal{B}(\chi) - 2\tau_0^2 \chi^2 \mathcal{B}(\chi)^2 + \mathcal{A}(\chi)^2, \\ \mathcal{A}(0) = \mathcal{B}(0) = 1. \end{cases} \quad (16)$$

The solution of these ODEs (which, see below, is easily found numerically) determines LL_∞ and its associated dynamics. There is a subtle sense in which this dynamics must be equivalent to that in LAD, and yet cannot be – on the one hand, (16) is defined by a fixed point condition matching it to LAD, but on the other hand LL_∞ is second-order in time derivatives and so must be free from pre-acceleration. While it is clearly not possible to analyse ‘pre-acceleration before the field turns on’ in our constant field setup, we can easily see that LL_∞ is causal; the general form (9) expresses the particle velocity as an integral over past, not future times. It appears, then, that LL_∞ may match LAD to all-orders in τ_0 , but misses non-perturbative effects in τ_0 (runaways and pre-acceleration), the recovery of which would likely require trans-series resummation [63]. We will compare to LAD below, but first we investigate the properties of LL_∞ in more detail.

The dynamics of LL_∞ has the same *qualitative* features as the dynamics of the resummed LL_n : \mathcal{A}, \mathcal{B} are monotonically decreasing and as $\mathcal{B} \leq 1$, the RR force is no stronger than that of LL_1 ; as \mathcal{B} is always positive, u^+ and hence χ are always driven to zero. The asymptotics of the solutions to LL_∞ and LL_1 are therefore the same. There are however *quantitative* differences at large χ , as we will now see.

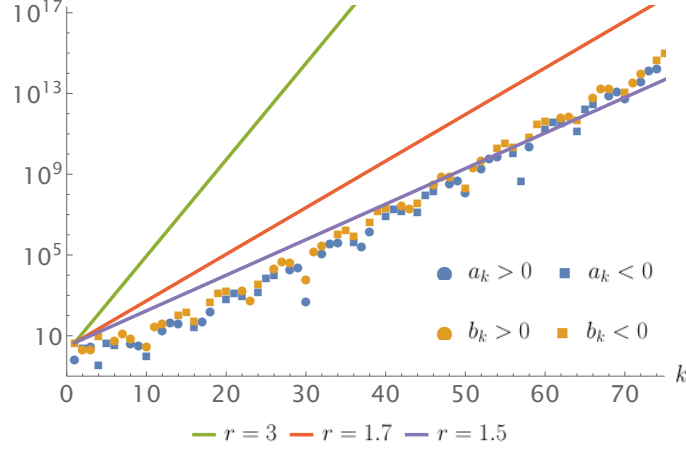


FIG. 4. The first few terms a_k (blue) and b_k (orange) of the series (17) and trendlines $|b_1|r^{k-1}$ (solid lines). Trend lines that lie above all data points give an estimate for the radius of convergence. The elements of the sequences have essentially random signs; to resolve these on a log scale, circles (squares) indicate positive (negative) values.

A. Large χ behaviour and convergence

We cannot solve the fixed-point ODEs analytically, but a series ansatz around $\chi = 0$ reproduces exactly the recursion relations (6). This prompts us to try a (Frobenius) series expansion for large χ ; we find

$$\mathcal{A}(\chi) = \sum_{k=0}^{\infty} a_k \left(\frac{3\tau_0}{\sqrt{2}} \chi \right)^{-(k+1)/2}, \quad \mathcal{B}(\chi) = \sum_{k=0}^{\infty} b_k \left(\frac{3\tau_0}{\sqrt{2}} \chi \right)^{-(k+3)/2}, \quad (17)$$

where the factor $3/\sqrt{2}$ has been introduced for convenience. The leading coefficients are $a_0 = 1$, $b_0 = 3$, with the higher orders determined by the recursion relation

$$\frac{1}{3} \begin{pmatrix} 3-k & 1 \\ -6 & 2-k \end{pmatrix} \begin{pmatrix} a_k \\ b_k \end{pmatrix} = \begin{pmatrix} -a_{k-1} + \frac{1}{9} \sum_{\ell} (\ell-3) a_{\ell} b_{k-\ell} \\ -b_{k-1} + \sum_{\ell} a_{\ell} a_{k-\ell} + \frac{1}{9} \sum_{\ell} (\ell-1) b_{\ell} b_{k-\ell} \end{pmatrix}, \quad (18)$$

the sums running over $1 \leq \ell \leq k-1$. Unlike the factorially growing A_k, B_k in (7), the coefficients a_k, b_k grow no faster than exponentially with k , as we illustrate for the first hundred in Fig. 4 and prove in Appendix B. The series (17) therefore have finite radii of convergence in $\chi^{-1/2}$; our proof provides a way to estimate the criterion for convergence as

$$\tau_0 \chi \gtrsim 1.37, \quad (19)$$

using the data in Fig. 4. We investigate the convergence of the series (17) in more detail in Fig. 5, by comparing against a numerical solution of the equations (16) for χ large; the latter requires setting conditions at $\chi = \infty$, where the fixed-point ODEs are singular, but since the series solution is convergent, we can set an initial condition $\mathcal{A}((\frac{3}{\sqrt{2}}\tau_0\chi)^{-1/2} = \varepsilon) = \varepsilon + a_1\varepsilon^2 + \dots$, thereby avoiding the singularity. Fig. 5 shows the partial sums indeed converging to the numerical solution when condition (19) holds. The numerical solution can be extended down to small values of χ , which can also be used to check on the resummation of the perturbative series, see below.

As the series (17) are convergent for large χ , we are encouraged to investigate the leading-order high- χ dynamics implied by LL $_{\infty}$ (16), which is encoded in the $k = 0$ terms of (17). These yield the large- χ asymptotics

$$\mathcal{A} \sim \frac{2^{1/4}}{3^{1/2}} \chi^{-1/2}, \quad \mathcal{B} \sim \frac{2^{3/4}}{3^{1/2}} \chi^{-3/2}, \quad (20)$$

in which we note that the powers appearing are the inverse of those typically associated with the Ritus-Narozhny conjecture [31]. As before we focus on the u^+ component, the equation for which, recall (8), here becomes

$$\int_{u_0^+}^{u^+} \frac{dv^+}{2\sqrt{v^+}} = - \int_0^{x^+} \left(\frac{m\mathcal{E}}{3\sqrt{2}\tau_0} \right)^{1/2} dy^+ \quad \implies \quad u^+(x^+) = u_0^+ \left(1 - \sqrt{\frac{m\mathcal{E}}{3u_0^+ \sqrt{2}\tau_0}} x^+ \right)^2. \quad (21)$$

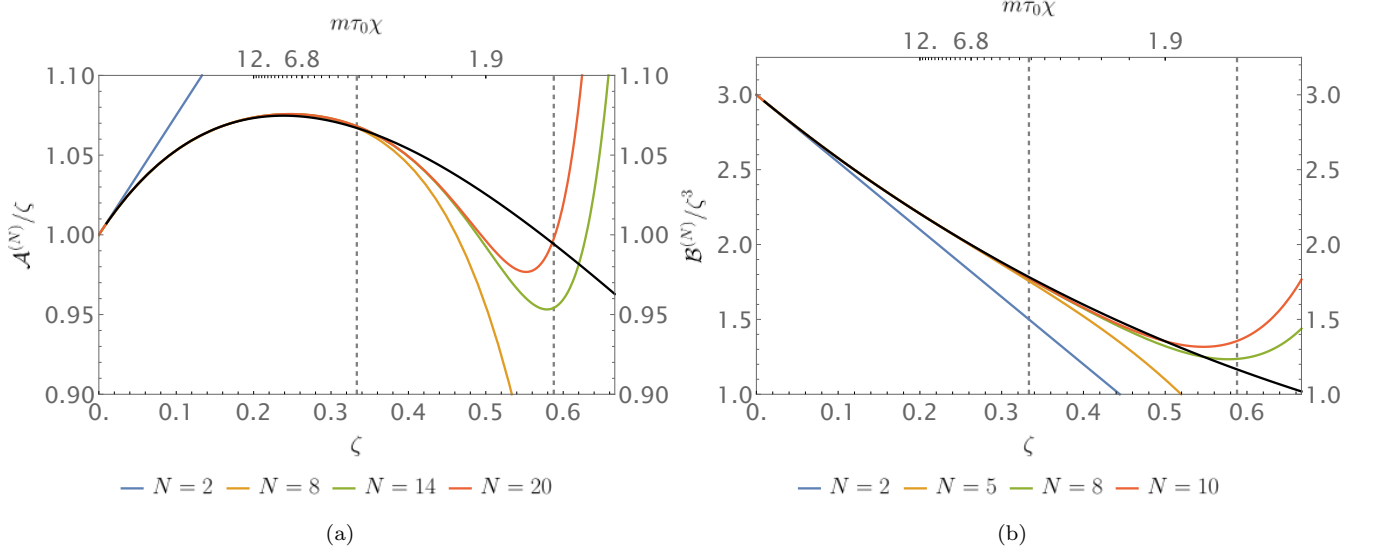


FIG. 5. Partial sums with N terms for the functions \mathcal{A}, \mathcal{B} (colour) and a numerical solution of the fixed-point ODE:s (black). The dashed lines indicate unconditional and estimated radii of convergence; see Appendix B. Here $\zeta = (3\tau_0\chi/\sqrt{2})^{-1/2}$.

Note, again, that the solution is explicitly causal: the solution $u(x^+)$ does not sample the field at future times. The solution shows that from an initial value of u_0^+ at $x^+ = 0$, the momentum u^+ drops to zero after a finite lightfront time, corresponding to the particle being accelerated to the speed of light [64]. This is qualitatively consistent with the general discussion of LL_∞ above, but the result (21) cannot be quantitatively valid for all times because as u^+ drops so does χ and we leave the high- χ region, hence the regime of validity of our approximation. We can estimate, directly from (21), that the particle remains at $\chi \gg 1$ only for

$$x^+ \lesssim \left(\frac{u_0^+ \tau_0}{m\mathcal{E}} \right)^{1/2}, \quad (22)$$

which can be understood as a measure of the time it takes for the initial lightfront energy mu^+ to be radiated away in a field of strength \mathcal{E} . Consequently the differences between LL_1 and LL_∞ will be most apparent if high χ is reached through high *energy*. Note that observables in QED have very different behaviours depending on whether χ is made large through high energy or high intensity, which has implications for the Ritus-Narozhny conjecture [65, 66].

IV. IMPROVED RESUMMATION

The preceding results demonstrate that we need the full χ -dependence of the functions \mathcal{A} and \mathcal{B} , even at high intensity where their series expansions are convergent, because the solution of LL_∞ is driven with lightfront time to the low- χ regime. In this section we use resummation to obtain an analytic approximation of the functions \mathcal{A} and \mathcal{B} which holds for $0 \leq \chi < \infty$.

We can see directly from (20) that the Borel-Padé resummed perturbative series considered in Sec. IIB, e.g. (14) and (15), do not give the correct high- χ scaling. This is due to Padé approximants always having integer power scaling, $N - M$, for large arguments. With integer powers, taking $N - M = 1$ is the closest we can get to the correct falloffs (20), which is why this choice leads to the most rapid convergence. We can improve upon this by instead using a resummation method that incorporates asymptotic data; one such method is “educated match” or Φ -Padé resummation [67]. In one variant of this method one constructs a Padé approximant

$$\Phi\text{P}_{\mathcal{A}}[N, M](t) = \frac{\sum_{\ell=0}^N c_\ell t^\ell}{1 + \sum_{j=1}^M d_j t^j} = \sum \frac{A_\ell t^\ell}{\ell!} \frac{\Gamma(\kappa)}{\Gamma(\kappa + \ell)} \frac{\Gamma(\lambda)}{\Gamma(\lambda + \ell)} + \mathcal{O}(\chi^{N+M+1}), \quad (23)$$

where the Borel transform has been replaced by a “hypergeometric transform” depending on two parameters κ and λ .

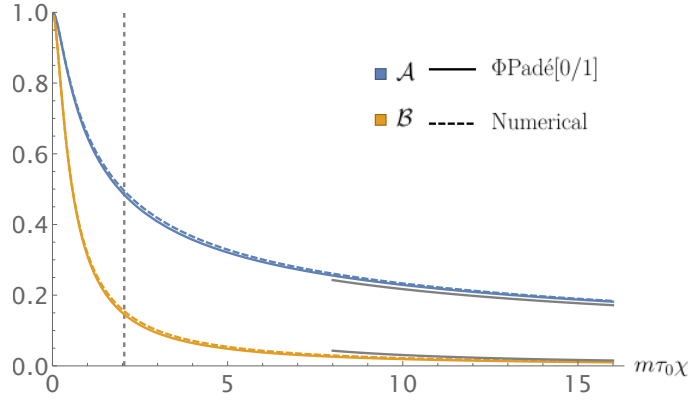


FIG. 6. Numerical solutions to the fixed-point ODEs (dashed) and the lowest-order Φ -Padé resumant (solid). Remarkable agreement is obtained using only a single perturbative coefficient, and by construction the correct asymptotic behaviour (grey curves) is reproduced. Vertical dashed grey lines indicate the radius of convergence of the large- χ series.

The resumant is obtained from the integral

$$\mathcal{A}(\chi) = \frac{1}{\Gamma(\kappa)\Gamma(\lambda)} \int_0^\infty dt e^{-t} t^{\kappa-1} \Phi P(t(\tau_0\chi)^2) U(1-\lambda, \kappa+\lambda+1, t), \quad (24)$$

where U is a confluent hypergeometric function [61, Ch. 13]. The asymptotic behaviour of the resumant is

$$\mathcal{A} \xrightarrow{\chi \rightarrow \infty} (\tau_0\chi)^{-2\lambda} \frac{\Gamma(\kappa-\lambda)}{\Gamma(\kappa)} + (\tau_0\chi)^{-2\kappa} \frac{\Gamma(\lambda-\kappa)}{\Gamma(\lambda)}, \quad (25)$$

from which κ, λ can be chosen to match a known asymptote. As with the order of the Padé approximant, different choices affect the rate of convergence. A convenient choice is $\lambda = 1$, as then one can make the replacement $U \mapsto 1$. The integral can then be performed analytically for the $[0, 1]$ approximants, superseding (14) and (15),

$$\mathcal{A}^{[0/1]}(\chi) = \frac{e^{\frac{1}{8\tau_0^2\chi^2}}}{8\tau_0^2\chi^2} E_{\frac{1}{4}}\left(\frac{1}{8\tau_0^2\chi^2}\right) \xrightarrow{\chi \rightarrow \infty} 0.73(\tau_0\chi)^{-1/2} \quad (26)$$

$$\mathcal{B}^{[0/1]}(\chi) = \frac{e^{\frac{1}{8\tau_0^2\chi^2}}}{8\tau_0^2\chi^2} E_{\frac{3}{4}}\left(\frac{1}{8\tau_0^2\chi^2}\right) \xrightarrow{\chi \rightarrow \infty} 0.76(\tau_0\chi)^{-3/2}, \quad (27)$$

where again E_n is an exponential integral [61, Ch. 5]. These have the correct asymptotic powers of χ by construction, with coefficients that are notably close to the correct values $2^{1/4}/\sqrt{3} \simeq 0.69$ and $2^{3/4}/\sqrt{3} \simeq 0.97$, despite us having used only a *single* perturbative coefficient to construct this lowest-order resumant. (Matching the coefficients exactly would require solving a transcendental equation for λ , leading to an integral we cannot perform analytically.) The analytic results (26) and (27) represent an excellent approximation to the numerically exact solution of (16) across the whole range of χ , as we demonstrate in Fig. 6. With this, we turn to a direct comparison of the dynamics implied by LL_∞ with those implied by LAD. (For comparisons with LL_1 , see Refs. [37, 68–70].)

V. BEYOND CONSTANT FIELDS: COMPARISON WITH LAD

Consider now LL_∞ for fields which are not constant, but which vary on time scales much longer than τ_0 . For such fields the derivative terms neglected in the derivation of LL_∞ are small. We show in Appendix A that our methods can be applied, and our conclusions remain robust, at leading order in derivatives. We therefore proceed to analyse motion in non-constant fields using a ‘locally constant field approximation’ [71], in which we simply promote f and χ appearing in LL_∞ to lightfront-time dependent variables.

Our test case is a particle created by some mechanism (e.g. the non-linear Breit-Wheeler process), in the peak of a circularly polarised pulse with a \sin^2 envelope and finite duration; the fieldstrength is $f_{\mu\nu} = n_\mu a'_\nu - a'_\mu n_\nu$, where

$$a'^\perp(\phi) = \mathcal{E} \sin^2 \frac{\phi}{4} \{\cos \phi, \sin \phi\}, \quad 0 \leq \phi \leq 4\pi, \quad \text{and zero otherwise.} \quad (28)$$

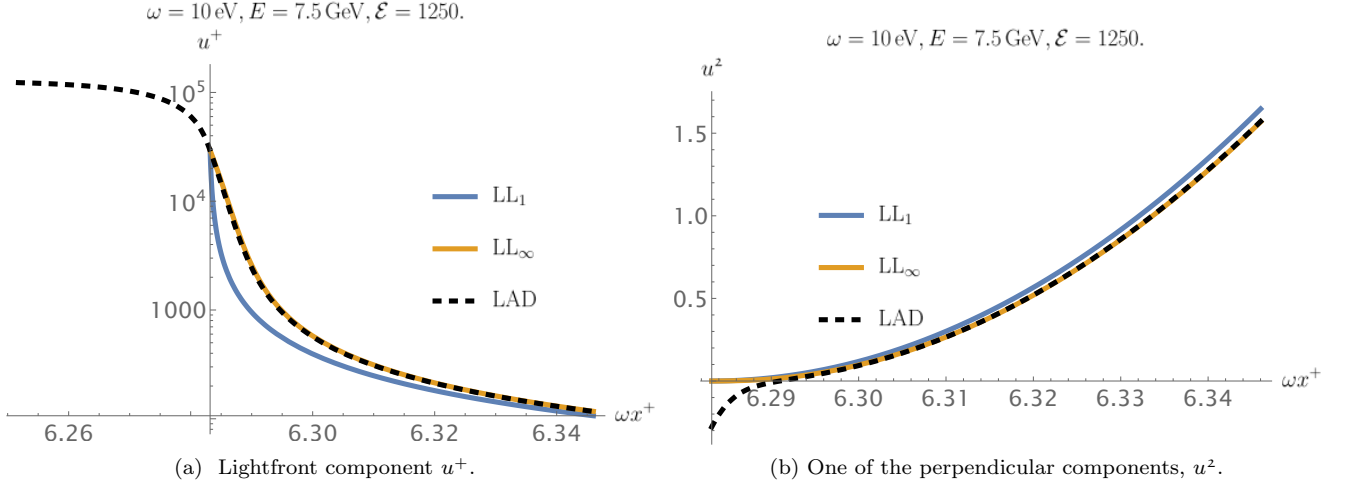


FIG. 7. Numerical solutions to LL_1 , LL_∞ , and LAD for a particle created at the peak of the pulse (28). The latter can be extended to earlier lightfront times with the field switched off, revealing its pre-accelerating nature. After the field has switched on, LL_∞ and LAD agree to within numerical precision. The dynamics of LL_1 differs significantly from these. In (b) the disagreement in transverse components near $\omega x^+ = 2\pi$ is due to numerical error, the scale of which is set by $u^+ \gg u^\perp$.

We consider this situation, rather than a particle already present before the pulse arrives, for the following reason. In a strong *pulsed* field, the large- χ regime where we expect differences between the equations of motion to become sizeable is never reached; even if the *peak* field and *initial* u^+ are taken to be large, almost all of the particle's lightfront energy will be radiated away long before it reaches the strong-field region [38, 44]. Our scenario thus mimics the situation where a charged particle encounters a 'step' at a given phase ϕ_i where the field is suddenly switched on.

In the differential formulation (1) of LAD, an additional boundary condition for \dot{u}^μ is needed, which can in principle be the physically motivated *final* condition $\dot{u}^+(\phi_f) = 0$, $\phi_f \geq 4\pi$, in our case, i.e. that the acceleration vanishes after the field has been turned off. 'Numerically exact' solutions of LAD must resolve the particle's proper time on the scale of τ_0 or less, as this multiplies the highest derivative in LAD. The analytical solution to LL_1 implies, though, that the proper time spent in the pulse is of order $3\pi^2\mathcal{E}^2\tau_0$, implying an unreasonable number of time steps for the case of interest, namely high intensities, $\mathcal{E}^2 \gg 1$. We will thus use the integro-differential [20, 23, 72, 73] formulation of LAD. It substitutes for the additional boundary condition a pre-accelerating integral over future times of the form

$$\iota \sim \int_{\tau}^{\infty} ds e^{-(s-\tau)/\tau_0} \mathcal{F}(s), \quad (29)$$

where \mathcal{F} is a certain function of u^μ and $f^{\mu\nu}$, see (C3) for details; ι must be known at $\phi_f = \omega x_f^+$ to solve backwards in time. The integral vanishes if the field vanishes for all future times, but again we need to resolve proper time on the scale τ_0 and so integrating backwards from $\phi = 4\pi$ is again unfeasible. However, after a very short time ε , almost all of the initial lightfront energy present at 'creation', $\phi_i = 2\pi$, will have been radiated away, and we can use the solution to LL_1 or LL_∞ to compute $\iota(\phi = 2\pi + \varepsilon)$, since by this point the particle has reached the regime where all the equations of motion agree. In any case, the solution is only weakly sensitive to the starting value of ι , as it is exponentially damped on the scale τ_0 . The details of our numerical discretisation scheme are further described in Appendix C.

We compare the solution to LAD thus obtained with that of LL_∞ in Fig. 7, using the analytical resummants (26) and (27) for \mathcal{A}, \mathcal{B} . Remarkable agreement is obtained in u^+ despite the minimum of information used in this representation of LL_∞ ; the error of at most $\simeq 5\%$ near $\phi = 2\pi$ is partly numerical, and otherwise stems from the relative error of the resummation (27) becoming significant for very large χ . (The much more apparent disagreement in u^\perp is predominantly error propagation from u^+ .) In Fig. 7 the low- χ regime is reached within a few percent of a cycle, well before the field changes significantly. This justifies *ex post* discarding derivative terms on this interval. We have also solved LAD and LL_∞ over one half of a cycle and seen that they still agree to within numerical precision.

As the field is only ever integrated in solving LAD, it is not a problem to explicitly turn off the field for $\phi \leq 2\pi$ and extend the LAD solution further back in time. The pre-acceleration contained in ι is thus explicitly seen in Fig. 7(a), becoming noticeable around $\Delta(\omega x^+) \simeq 0.01$ before the field switches on. Estimating $\Delta\tau/\tau_0 \approx \Delta(\omega x^+)/\omega\tau_0 u^+(2\pi) \approx 3.5$, the pre-acceleration occurs, as expected, over a few τ_0 of proper time.

We mention finally that the accuracy of the approach used here can be improved further by constructing piecewise approximations to \mathcal{A} and \mathcal{B} ("stitching"), in which one uses the resummed expressions below some cutoff in χ , and the convergent large- χ expansion above the cutoff. As few as 5 terms each for Φ -Padé resummation and series expansion

suffice to approximate \mathcal{A} and \mathcal{B} to within a few percent over the full range of χ , leading to a relative error in u^+ of $\simeq 1\%$. This confirms that LL_∞ reproduces the physical solution of LAD, while being explicitly causal and second-order in time.

VI. CONCLUSIONS

In this paper we have applied iterated reduction of order to the LAD equation of motion in classical electrodynamics. This generates an infinite sequence of possible approximations to LAD which is typically truncated at leading order, yielding the well-known Landau-Lifshitz equation. Working for simplicity with a constant crossed field background, we have investigated the ultimate fixed point of the iteration, the equation LL_∞ . (The constant field approximation can be viewed as the leading order in a derivative expansion, with the extension to next-to-leading order sketched in Appendix A.)

We were able to map this problem to a nonlinear initial value problem for two functions \mathcal{A}, \mathcal{B} of the invariant χ , which we have approached through (i) a direct numerical solution, which provides an ‘exact’ benchmark, (ii) a perturbative, weak-field analysis, (iii) resummation of the resulting asymptotic weak-field expansion, (iv) a strong-field expansion, which is convergent.

Regarding (ii), we can understand why the weak-field perturbative expansion is divergent through a variation of Dyson’s argument [74] that the perturbation series of QED has zero radius of convergence due to the instability of the vacuum against pair production when $\alpha < 0$. We adapt this for our case as follows: since $\alpha > 0$, an accelerating particle emits radiation in the direction of its motion, and thus RR opposes the motion. If α were negative, though, particles would instead *gain* momentum due to the reaction, enhancing the emission of radiation. In such a universe particles in magnetic or Coulomb fields would spiral out, not in. This instability cannot be connected to the physical universe, with $\alpha > 0$, by a perturbative expansion around $\alpha = 0$.

Turning to (iii), Borel-Padé resummation gives a good approximation to the numerically exact solution. Using improved resummation methods, only a single perturbative coefficient, along with data from the convergent asymptotic series (iv), is needed to yield a resummant which agrees to within a few percent for all χ . We thus have an approximate analytic expression for the fixed-point, all-orders LL equation, LL_∞ . It is second order in time derivatives, and can be solved numerically to arbitrary precision. An important qualitative feature of LL_∞ is that (because the function \mathcal{B} is always positive) χ is always driven to 0. As a result large-time dynamics are governed by LL_1 , and asymptotic quantities match those of LL_1 . (In particular the motion becomes aligned with the ‘radiation-free direction’ in which there is no acceleration transverse to the motion, minimising radiation losses, as also holds for LAD [44, 48, 51, 75].) Without resummation, though, it follows that the practical, optimal, order of truncation for the series generated by iterated reduction of order is 1.

Interestingly, though, we note that the short-time dynamics of LL_∞ depends on whether it is, in the composite parameter $\chi = u^+ \mathcal{E}$, the energy u^+ or the intensity \mathcal{E} which is made large. This is also the case in QED in strong fields, where the high-energy and high-intensity limits of observables are drastically different, scaling with powers or logarithms, respectively [65, 66].

Continuing with connections to the quantum theory, a common adage is that LL_1 overpredicts the strength of RR compared to QED. In this sense LL_∞ gives results ‘closer’ to QED: because the functions \mathcal{A}, \mathcal{B} are strictly decreasing, with a maximum of 1, LL_∞ predicts less RR than LL_1 . This is explicitly seen in Fig. 7.

Both LL_1 and LAD are known to be consistent with QED to leading order in α [76–78], with a very recent resummation of quantum RR in plane waves [29] recovering LL_1 to all orders in α , but leading order in the pulse duration and intensity. As this is a result for long times and in the high-*intensity* limit, it is still consistent with LL_∞ , which differs markedly from LL_1 only for short times, except in the high-*energy* limit. Making more precise, quantitative statements about the relation between LL_∞ , LAD, and QED requires going beyond leading order in the pulse duration or understanding the high-energy limit of strong-field QED, which remains challenging.

ACKNOWLEDGMENTS

The day after this manuscript was uploaded to the arXiv, Ref. [79] appeared, Sec. IV of which also treats resummation in the context of LAD. Where we have resummed on the level of the equation, Ref. [79] has resummed on the level of the solution. We have checked that the approaches are consistent, yielding e.g. the same acceleration at $x^+ = 0$, as the coefficients in (60) and (61) of Ref. [79] match our (7). We thank Greger Torggrímsson for in-depth discussions on this.

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Appendix A: LL_∞ to leading order in derivatives

The characterisation of LL_∞ as a fixed-point of reduction of order is not limited to a constant field, which serves only to limit the tensor that can appear. In a non-constant field LL_∞ must contain, among others, terms with arbitrary number of derivatives,

$$\dot{u}^\mu = \sum_{\ell=0} \mathcal{A}_\ell (u \cdot \partial)^\ell f^{\mu\nu} u_\nu + \mathcal{B}_\ell [\mathcal{P}(u \cdot \partial)^\ell f^2]^{\mu\nu} u_\nu + \dots \quad (\text{A1})$$

where the functions $\mathcal{A}_\ell, \mathcal{B}_\ell$ depend on not just the one invariant y , but also on, among others, $y_\ell \propto [(n \cdot \partial)^\ell f_{\mu\nu} u^\nu]^2$. The fixed-point condition is then, in general, an infinite tower of PDE:s in an infinite-dimensional space.

Truncating the expansion at $\ell = 0$ is justified as long as $\omega\tau_0$ remains small, where ω is a typical frequency of the field, and the $\mathcal{A}_\ell, \mathcal{B}_\ell$ don't blow up. We cannot prove this in full generality, but we can consider the next simplest case in the derivative expansion.

To first order in derivatives LL_∞ must have the form

$$\dot{u}_\lambda = (A_1 + zB_1)f^{\mu\nu}u_\nu + \frac{1}{\tau_0}(yA_2 + zB_2)\left(\frac{n^\mu}{n \cdot u} - u^\mu\right) + \tau_0 A_3 f^{\mu\nu,\rho}u_\nu u_\rho \quad (\text{A2})$$

where the invariant z is

$$z = \tau_0^3 u_\mu f^{\mu\nu} f_{\nu\rho,\sigma} u^\rho u^\sigma = \tau_0^3 (n \cdot u)^3 (a'^\perp \cdot a''^\perp) \quad (\text{A3})$$

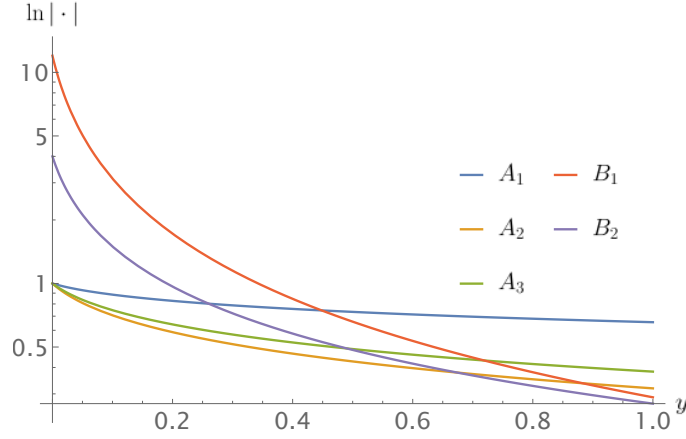


FIG. 8. Numerical solutions for the coefficient functions A_i in the leading-order derivative expansion of LL_∞ , (A4)–(A5).

and the coefficient functions A_i, B_i depend only on χ . The initial conditions are set by LL_1 as $A_3(y=0) = 1$ and a simple calculation to find $B_1(0) = -12, B_2 = 4$.

Applying reduction of order, identifying terms by tensor structure and order by order in z , the fixed-point condition now results in the initial value problem

$$\begin{aligned}
 -2y^2 A_2 A_1' - 2y A_1 A_2 + 1 &= A_1 \\
 -2y^2 A_2 A_2' - 2y A_2^2 + A_1^2 &= A_2 \\
 A_1 - 2y^2 A_2 A_3' - 3y A_2 A_3 &= A_3 \\
 2(1 - y B_2) A_1' - 2y^2 A_2 B_1' - 5y A_2 B_1 - 2A_1 B_2 &= B_1 \\
 2(1 - y B_2) y A_2' + 2A_2 - 2y^2 A_2 B_2' - 5y A_2 B_2 + 2y A_1 B_1 + 2A_1 A_3 &= B_2
 \end{aligned} \tag{A4}$$

$$A_1(0) = A_2(0) = A_3(0) = 1 \quad B_1(0) = -12 \quad B_2(0) = 4. \tag{A5}$$

The first two equations and initial conditions are, naturally, precisely the fixed-point conditions at zeroth order in derivatives. This system can be solved numerically, Fig. 8, showing that all four functions fall off monotonically with y . We can determine their asymptotics for $y \gg 1$ by assuming a power law falloff, resulting in

$$A_1 \sim \frac{2^{1/4}}{3} y^{-1/4} \quad A_2 \sim \frac{2^{3/4}}{\sqrt{3}} y^{-3/4} \quad A_3 \sim \frac{1}{2^{3/2}} y^{-1/2} \quad B_1 \sim -\frac{1}{2^{3/4}} y^{-3/2} \quad B_2 \sim \frac{1}{3} y^{-1}. \tag{A6}$$

Furthermore, the asymptotic forms suggest that A_3, A_4 also have convergent series expansion around $\chi = \infty$. As the structure of the initial-value problem is the same as without derivatives, a recursion relation similar to (18) can be found, and the proof in Appendix B will apply.

Appendix B: Proof of convergence of the series around $\chi = \infty$

We will prove that the series (18) has a finite radius of convergence. First, assume that k is large enough that we can disregard terms subleading in k , viz.,

$$\begin{pmatrix} a_k \\ b_k \end{pmatrix} \approx -\frac{1}{3k} \begin{pmatrix} \sum \ell a_\ell b_{k-\ell} \\ \sum \ell b_\ell b_{k-\ell} \end{pmatrix} + \mathcal{O}(1/k) \tag{B1}$$

Let β_k be the sequence

$$\beta_k = \frac{1}{2\lambda k} \sum_{\ell=1}^{k-1} \ell \beta_\ell \beta_{k-\ell} = \frac{1}{4\lambda} \sum_{\ell=1}^{k-1} \beta_\ell \beta_{k-\ell} \quad k > 1. \tag{B2}$$

If $|\beta_\ell| > |b_\ell|, |a_\ell|$ for $\ell < k$, the derivative term is bounded by β_k , so a_k, b_k differ by at most $\mathcal{O}(1/k)$. Thus β_k will bound $|a_k|, |b_k|$ as long as it does so until k becomes sufficiently large. We will now prove that through induction on k that

$$\beta_k \leq \lambda \frac{\beta_1^k}{\lambda^k}; \quad (\text{B3})$$

the base case $k = 1$ is obviously true. To reduce clutter let's introduce the shorthand

$$[n_1, n_2, \dots, n_N] := \prod_{i=1}^N \beta_{n_i} \quad (\text{B4})$$

First note that $\beta_2 = \beta_1^2/(4\lambda)$; then for $k > 2$ write

$$\beta_k = \frac{\beta_1}{2\lambda} \beta_{k-1} + \frac{1}{4\lambda} \sum_{n_i \geq 2} [n_1, n_2] \delta(k - \sum n_i) \quad (\text{B5})$$

$$= \frac{\beta_1}{2\lambda} \beta_{k-1} + \frac{1}{(4\lambda)^2} \sum_{n_i \geq 2} [\ell_1, \ell_2, n_1] \delta(k - \ell_1 - \ell_2 - \sum n_i) \quad (\text{B6})$$

“splitting” either of the summation indices using the definition (B2). Since $n_1 \geq 2$ the maximally indexed factors that appear in the triple product are β_{k-3} and β_{k-4} , according to the minimal sums of ℓ_i . The former has combinatorial weight 1 and the latter 2. Thus,

$$\beta_k = \frac{\beta_1}{2\lambda} \beta_{k-1} + \frac{\beta_1^2}{8\lambda^2} \beta_{k-2} + \frac{\beta_1^3}{16\lambda^3} \beta_{k-3} + \frac{1}{(4\lambda)^3} \sum_{n_i \geq 2} [n_1, n_2, n_3] \delta(k - \sum n_i) \quad (\text{B7})$$

and the triple product can be split into a quadruple product with maximal indices $k-4$ and $k-5$. The combinatorial weights are 1×2 and 2×2 , respectively. In general, splitting an index in a product with N indices (which will have a factor $(4\lambda)^{-N+1}$ from previous splits, will work out like

$$\sum_{n_i \geq 2} [n_1, \dots, n_N] \delta(k - \sum n_i) = \frac{1}{4\lambda} \sum_{n_i \geq 2} [\ell_1, \ell_2, n_1, \dots, n_{N-1}] \delta(k - \ell_1 - \ell_2 - \sum n_i) \quad (\text{B8})$$

$$= \frac{N-1}{4\lambda} \beta_2^{N-2} (\beta_1^2 \beta_{k-2N+2} + 2\beta_1 \beta_2 \beta_{k-2N-1}) + \frac{1}{4\lambda} \sum_{n_i} [n_1, \dots, n_{N+1}] \delta(k - \sum n_i) \quad (\text{B9})$$

$$\leq \frac{\beta_1^{2N-2}}{2^{N-1} \lambda^{N-1}} \beta_{k-2N+2} + \frac{\beta_1^{2N-1}}{2^{N-2} \lambda^N} \beta_{k-2N-1} + \frac{1}{4\lambda} \sum_{n_i} [n_1, \dots, n_{N+1}] \delta(k - \sum n_i) \quad (\text{B10})$$

Repeating and using the induction hypothesis (B3) we can bound β_k by a geometric series:

$$\beta_k \leq \sum_{\ell=1}^k \frac{\beta_1^\ell}{(2\lambda)^\ell} \beta_{k-\ell} \leq \lambda \frac{\beta_1^k}{\lambda^k} \sum_{\ell=1}^k \frac{1}{2^\ell} = \beta_1 \frac{\beta_1^k}{\lambda^k}. \quad (\text{B11})$$

To get a bound for our b_k :s we should take $\beta_1 = |b_1| = \frac{9}{2}$ and $\lambda \leq 3/2$ which implies $\rho \geq \frac{3}{2|b_1|} = 1/3$. As evidenced in Fig. 5, the actual radius of convergence is larger. This is because a_ℓ, b_ℓ are more or less alternating. To take this into account, we can consider the analogous sequence $\tilde{\beta}_k$ where we set $\tilde{\beta}_2 = -\tilde{\beta}_1^2/4\lambda$. Now realise that every time we pick an $\ell_i = 2$, we get a minus sign from this, so the geometric series (B11) becomes alternating, and we can actually bound

$$|\tilde{\beta}_k| \leq \lambda \left(\frac{|\tilde{\beta}_1|}{2\lambda} \right)^k. \quad (\text{B12})$$

Thus, our sequences a_k, b_k are bounded by

$$\left| \frac{a_k}{a_1} \right|, \left| \frac{b_k}{b_1} \right| \leq r^k \quad (\text{B13})$$

as long as r is large enough to both satisfy $r > \frac{1}{2} \max(|a_1|, |b_1|)$ and provide a bound even for $1 \ll k$. We can estimate such an r by a linear fit to a few hundred elements, or by graphical estimation; see again Fig. 4: if r is such that all the points are below its trendline, the radius of convergence in $(\frac{3\tau_0}{\sqrt{2}}\chi)^{-1/2}$ is at least $1/r$. We see that $r = \frac{3}{2}$ fails to give a bound for the a_k, b_k around $k \approx 50$, but $r \approx 1.7$ – found by fitting – provides the bound (19), $m\tau_0\chi \gtrsim 1.37$.

Appendix C: Numerical scheme for LAD

The runaway solutions can be removed by writing LAD in an integro-differential form; that presented in Ref. [73] is suitable. It is obtained by rearranging LAD into

$$f^{\mu\rho}u_\rho + \tau_0 u^\mu \dot{u}^\rho \dot{u}_\rho = -\tau_0 e^{\tau/\tau_0} \frac{d}{d\tau} \left(e^{-\tau/\tau_0} \dot{u}^\mu \right) \quad (C1)$$

tensoring by u^ν and anti-symmetrising to kill the term quadratic in \dot{u}^μ , viz.,

$$-\frac{e^{-\tau/\tau_0}}{\tau_0} (f^{\mu\rho}u_\rho u^\nu - f^{\nu\rho}u_\rho u^\mu) = u^{[\nu} \frac{d}{d\tau} \left(e^{-\tau/\tau_0} \dot{u}^{\mu]} \right) = \frac{d}{d\tau} \left(e^{-\tau/\tau_0} u^{[\nu} \dot{u}^{\mu]} \right). \quad (C2)$$

The equation is now solved by formal integration ($dx^+ = x^+ d\tau$) and dotting with $u^\nu(x^+)$,

$$\frac{du^\mu}{dx^+} = \frac{e^{\tau/\tau_0}}{\tau_0} \frac{u_\nu}{u^+} \int_{x^+}^\infty \frac{dy e^{-\tau(y)/\tau_0}}{u^+(y)} [f^{\mu\rho}u_\rho u^\nu - f^{\nu\rho}u_\rho u^\mu](y), \quad (C3)$$

free of runaways as long as $f^{\mu\nu}$ vanishes sufficiently quickly.

Our discretisation scheme is a simple finite difference scheme. Letting $\tilde{\tau} = \tau/\tau_0$, and subscript k indicating evaluation at ϕ_k , we proceed according to

$$u_{k+1}^\mu - u_k^\mu = \Delta\phi \frac{1}{\tau_0\omega} \frac{u_k^\nu}{u_k^+} \int_\phi^\infty d\phi' \frac{e^{(\tilde{\tau}(\phi) - \tilde{\tau}(\phi'))}}{u^+(\phi')} [\tilde{f}^{\mu\rho}u_\rho u_\nu - \tilde{f}^{\nu\rho}u_\rho u^\mu](\phi') \quad \tilde{f}^{\mu\nu} = n^{[\mu} \frac{da^{\nu]}}{d\phi} \quad (C4)$$

where we evaluate the integral by the trapezoid method,

$$I_k = \int_{\phi_k}^\infty e^{\tilde{\tau}_k - \tilde{\tau}(y)} \dots dy = \int_{\phi_k}^{\phi_{k-1}} e^{\tilde{\tau}_k - \tilde{\tau}(y)} \dots dy + e^{\tilde{\tau}_k - \tilde{\tau}_{k-1}} \int_{\phi_{k-1}}^\infty e^{\tilde{\tau}_{k-1} - \tilde{\tau}(y)} \dots dy \quad (C5)$$

$$\approx \frac{\Delta\phi}{2} (f(\phi_k, u_k) + e^{\tilde{\tau}_k - \tilde{\tau}_{k-1}} f(\phi_{k-1}, u_{k-1})) + e^{\tilde{\tau}_k - \tilde{\tau}_{k-1}} I_{k-1}. \quad (C6)$$

The trapezoid method is also used to evaluate the proper time step,

$$\tilde{\tau}_k - \tilde{\tau}_{k-1} = \frac{\tau_k - \tau_{k-1}}{\tau_0} = -\frac{\Delta\phi}{2\omega\tau_0} \left(\frac{1}{u_k^+} + \frac{1}{u_{k-1}^+} \right); \quad (C7)$$

while $I_0 = \int_{\phi_0}^\infty dy \cdot$ is evaluated using a solution of LL₁, since by assumption all later times are in the low- χ regime where LL₁ and LAD agree; this initialisation decays over a very short phase interval, as u^+ is small near ϕ_0 , and the solution is not sensitive to it. While the equation is written in covariant form, we treat only u^+ , u^\perp as independent, fixing the final component through the mass-shell condition $u^- = \frac{1+(u^\perp)^2}{u^+}$.