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ACTION-AT-A-DISTANCE FORMULATION OF FIELD THEORIES FOR ATOMIC NUCLEI

H. FELDMEIER

(Invited talk presented at 'Encuentros Relativistas Espanoles 93', Salas, Asturias, 7-10 Sept. 1993)

Gesellschaft für Schwerionenforschung mbH
Postfach 11 05 52 · D-64220 Darmstadt · Germany

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Hans Feldmeier

Gesellschaft für Schwerionenforschung mbH,
Postfach 110 552, D-64220 Darmstadt, Germany

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1. Relativistic mean field theory for nuclei and nuclear matter

In effective theories for hadronic systems the interaction between the nucleons is mediated by massive mesons [1] while the quark and gluon degrees of freedom are not resolved. The saturation in the binding energy and the strong spin–orbit splitting observed in nuclei indicate that a relativistic description with strong Lorentz–scalar and Lorentz–vector fields for the attraction and repulsion, respectively, might be useful, even though the Fermi velocity is only one third of the speed of light. In the most simple relativistic model, the so called Walecka model, the mesons are described by c-number mean fields [2]. In this and related models the long range attraction is due a scalar field $\phi(x)$ with a boson mass around $\mu_S = 500$ MeV and the short range repulsive part is represented by a vector field $A^\alpha(x)$ with mass $\mu_V = 750$ MeV. One usually starts from a lagrangian of the following type

$$\begin{aligned}\mathcal{L}(x) &= \bar\psi(x)(\gamma^\alpha i\partial_\alpha - M)\psi(x) + g_S\bar\psi(x)\psi(x)\phi(x) \\ &- g_V\bar\psi(x)\gamma^\alpha\psi(x)A_\alpha(x) \\ &- \frac{1}{2}\phi(x)(\partial_\alpha\partial^\alpha + \mu_S^2)\phi(x) - \mathcal{U}(\phi(x)) \\ &+ \frac{1}{2}A^\alpha(x)(\partial_\beta\partial^\beta + \mu_V^2)A_\alpha(x)\end{aligned} \quad (1)$$

Self interaction terms $\mathcal{U}(\phi(x))$ are often added to increase the effective mass and improve surface properties of nuclei. The corresponding field equations are

$$\Big\{(i\partial_\alpha - g_V A_\alpha(x)) - (M - g_S\phi(x))\Big\}\psi(x) = 0 \quad (2)$$

for the nucleons and

$$(\partial^\beta\partial_\beta + \mu_S^2)\phi(x) = g_S\,\bar\psi(x)\psi(x) - \frac{d\mathcal{U}}{d\phi}(\phi(x)) \quad (3)$$

for the scalar field and

$$(\partial^\beta\partial_\beta + \mu_V^2)A^\alpha(x) = g_V\,\bar\psi(x)\gamma^\alpha\psi(x) \quad (4)$$

for the vector field. In the mean–field approximation the source terms are replaced by their expectation values. The scalar density $\langle\bar\psi(x)\psi(x)\rangle$ and vector current density $\langle\bar\psi(x)\gamma^\alpha\psi(x)\rangle$ are calculated by summing over occupied positive energy states.

In a nucleus like ^{208}Pb the mean fields are of the order of $g_V A^0 \approx -g_S\,\phi \approx \frac{1}{3}M$ which explains the strong spin–orbit splitting. On the other hand the classical potential energy which is the difference of both, $g_V A^0 - g_S\,\phi$, is only about -50 MeV, or $\frac{1}{20}M$. One has to realize that the fields $\phi(x)$ and $A^\mu(x)$ do not represent quantized fields of existing mesons. They are rather effective fields with adjustable coupling strengths and include many effects like exchange terms, renormalization or double pion exchange. The lightest particle, the pion, does not occur as a mean field due to its negative parity.

With only a few parameters these models are very successful in reproducing properties of many nuclei in or close to the ground state, see for example ref. [3]. But also hot and dense nuclear or neutron matter has been described within this relativistic frame work e.g. in ref. [4]. It is therefore near at hand to use the same relativistic model for the description of nucleus–nucleus collisions.

2. Relativistic time–dependent mean fields for colliding nuclei

In collisions the time–dependence of the fields adds the new aspect of radiation. Bai et al. [5] and Weber et al. [6] have solved numerically the wave equations (3) and (4) in the mean–field approximation together with equations for the scalar density and vector current. These calculations are very time consuming and can hardly be repeated for other cases. Within the numerical inaccuracies their solutions (for collisions with 1 GeV per nucleon) show that the radiated fields are rather weak. The reason is that the fields are massive and radiation requires Fourier frequencies in the source densities which are at least about the meson mass.

More serious than the technical problem is the conceptual problem how to interpret these radiated c–number fields. There is the narrow ω–resonance at $m_\omega = 782$ MeV which can be regarded as the main contribution to the $A^\alpha(x)$ field, but there is no such resonance which corresponds to the $\phi(x)$ field of a scalar meson with mass around 500 MeV. Thus, the physical interpretation of freely travelling ϕ fields or field quanta is obscure.

Therefore I proposed [7, 8] to exclude radiation of ficticious mesons from the very beginning by using the action–at–a–distance formulation with the symmetric Green's function of Wheeler and Feynman [9].

$$G_{S,A}(x-y) = \frac{1}{2}\left(G_{S,A}^{advanced}(x-y) + G_{S,A}^{retarded}(x-y)\right) \tag{5}$$

The formal solution of the wave equation (3) (for $\mathcal{U}(\phi)=0$)

$$\phi(x) = g_S \int d^4y \; G_S(x-y)\langle\bar{\psi}(y)\psi(y)\rangle \tag{6}$$

and of eq. (4)

$$A^\alpha(x) = g_V \int d^4y \; G_V(x-y)\langle\bar{\psi}(y)\gamma^\alpha\psi(y)\rangle \tag{7}$$

fulfills the desired boundary condition of vanishing incoming and outgoing free fields.

Eliminating the fields from the lagrangian (1) leads to the lon–local action which contains only nucleon variables:

$$\begin{aligned}\int d^4x \; \mathcal{L}(x) &= \int d^4x \; \bar{\psi}(x)(\gamma^\alpha i\partial_\alpha - M)\psi(x) \\ &+ \frac{1}{2}g_S^2 \int d^4x \; d^4y \; \bar{\psi}(x)\psi(x)\, G_S(x-y)\, \bar{\psi}(y)\psi(y) \\ &- \frac{1}{2}g_V^2 \int d^4x \; d^4y \; \bar{\psi}(x)\gamma^\alpha\psi(x)\, G_V(x-y)\, \bar{\psi}(y)\gamma_\alpha\psi(y)\end{aligned} \tag{8}$$

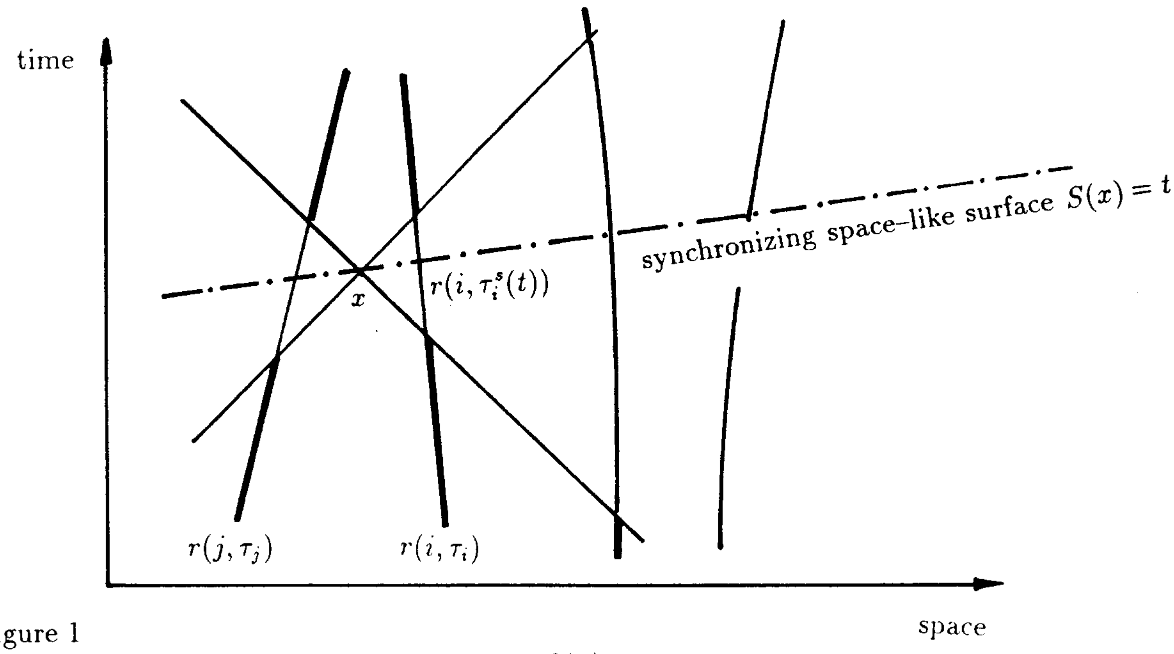

Figure 1
World lines and synchronizing hypersurface $S(x)$

Inside a nucleus the nucleons cannot be localized well enough to treat them as classical particles. But their phase-space distribution can be represented by means of test particles [7]. For example the scalar density and vector current are represented as

$$\rho_S(x) := \langle\bar{\psi}(x)\psi(x)\rangle \quad \rightarrow \quad \frac{1}{N}\sum_{j=1}^{A\cdot N}\int d\tau_j \; \delta^4(x - r(j,\tau_j)) \tag{9}$$

$$j^\mu(x) := \langle\bar{\psi}(x)\gamma^\mu\psi(x)\rangle \quad \rightarrow \quad \frac{1}{N}\sum_{j=1}^{A\cdot N}\int d\tau_j \; \delta^4(x - r(j,\tau_j))\, u^\mu(j,\tau_j) \;, \tag{10}$$

where $r(j, \tau_j)$ denotes the world line of test particle j at its proper time τ_j and $u(j, \tau_j)$ its four-velocity.

In this representation the scalar field at a space-time point x is given by integrals over past and future proper times τ_j as

$$\phi(x) = \hat{g}_S \sum_{j=1}^{A\cdot N}\int d\tau_j \; G(x - r(j,\tau_j)) \;. \tag{11}$$

N is the number of test particles per nucleon and A is the number of nucleons. $\hat{g}_S = g_S/N$ is the proportionate charge of each test particle.

If the test particles (or pieces of the phase-space density) are not too strongly accelerated one may expand each word line around the proper time τ_j^s which is determined from $x^0 = r^0(j, \tau_j^s)$ as

$$r^\alpha(j,\tau_j) = r^\alpha(j,\tau_j^s) + (\tau_j - \tau_j^s)u^\alpha(j,\tau_j^s) + \frac{1}{2}(\tau_j - \tau_j^s)^2\, a^\alpha(j,\tau_j^s) + \cdots \;. \tag{12}$$

Neglection of the quadratic term with the acceleration $a^\alpha(j, \tau_j^s)$ and all higher powers will be called "small acceleration approximation" and leaves us with a straight world line from the past to the future.

The small acceleration approximation is best in the vicinity of $r^\alpha(j,\tau_j^s)$ and becomes worse further away. A world line which hits the light cone, centered at x, far away from x may be badly approximated by eq. (12), but for short range interactions a distant particle does not contribute anymore to the field at x, c.f. fig. 1.

Thus, the first condition for the approximation is that the range μ^{-1} is small compared to the curvature of the world lines, i.e. the inverse of the acceleration. The second condition is weak radiation, which is fulfilled when the acceleration is small compared to the meson mass μ. Both conditions are actually the same, namely

$$\mid a_\mu a^\mu \mid \ll \mu^2. \tag{13}$$

Here one sees the difference to electrodynamics where the action–at–a–distance formulation cannot be used since $\mu = 0$. The range of the Coulomb interaction is infinite and even small accelerations lead to radiation.

In nuclear physics the assumption of small accelerations is justified if the ϕ and A^α fields are only meant to be the mean field part of the nucleon–nucleon interaction in a hadronic surrounding. The hard collisions between individual nucleons which are due to the repulsive core will cause large accelerations and also create new particles. These hard collisions cannot be described within the Walecka model, therefore it is consistent to regard $\phi(x)$ and $A^\alpha(x)$ as Hartree mean–fields which bring about only small accelerations and which are not radiated away from their sources.

Inserting the straight world line into eq. (11) results in the easily understood situation that the field at a point x is just the sum of Lorentz–boosted Yukawa potentials which are travelling along with the charges:

$$\phi(x) = \frac{\hat{g}_S}{4\pi} \sum_j \frac{\exp\left\{-\mu_S\sqrt{-R(x,j)^2}\right\}}{\sqrt{-R(x,j)^2}} \, , \tag{14}$$

where $R(x,j)^2$ is given by

$$R(x,j)^2 = (x - r(j,\tau_j^s))^2 - \left[(x^\alpha - r^\alpha(j,\tau_j^s))\, u_\alpha(j,\tau_j^s)\right]^2 \tag{15}$$

The metric is chosen as $(1,-1,-1,-1)$. The vector field is derived in an analogue fashion as

$$A^\alpha(x) = \frac{\hat{g}_V}{4\pi} \sum_j \frac{\exp\left\{-\mu_V\sqrt{-R(x,j)^2}\right\}}{\sqrt{-R(x,j)^2}} \, u^\alpha(j,\tau_j^s) \, . \tag{16}$$

At this level of the approximation the causality problem with the advanced part of the Green function is not present because the retarded and the advanced fields are identically the same when they are created by charges moving on straight world lines. Therefore one may regard the fields as retarded only.

Variation of the action (21), which is given in the following section, yields the equations of motion of the test particles.

$$\begin{aligned} \frac{d}{d\tau} r^\alpha(i,\tau) &= u^\alpha(i,\tau) \\ \frac{d}{d\tau} u^\alpha(i,\tau) &\equiv a^\alpha(i,\tau) \end{aligned}$$

$$= -\frac{1}{\widehat{M} - \hat{g}_S\phi(i,\tau)}\Big[\partial^\alpha\phi(i,\tau) - (\partial_\beta\phi(i,\tau)u^\beta(i,\tau))\,u_\alpha(i,\tau)$$
$$+(\partial^\alpha A^\beta(i,\tau) - \partial^\beta A^\alpha(i,\tau))u_\beta(i,\tau)\Big]\,. \tag{17}$$

The following short hand notation has been used:

$$\phi(i,\tau) = \phi(r_i,\tau) \quad \text{and} \quad \partial^\alpha\phi(i,\tau) = \frac{\partial}{\partial r_\alpha(i,\tau)}\phi(r_i,\tau) \tag{18}$$

and analogously for $A^\alpha(i,\tau)$. If one inserts the approximated fields of eqs. (14) and (16) and their derivatives at the location $x = r(i,\tau)$ one obtains the equations of motion for the test particles in the small acceleration approximation [7]. The summation in eqs. (14) and (16) exclude of course test particle i. The synchronization is realized by choosing the proper times τ_j^s of all particles which create the fields felt by particle i such that $r^0(j,\tau_j^s) = r^0(i,\tau)$ for all j in a selected coordinate frame.

These equations of motion have already been known in Predictive Relativistic Mechanics (PRM) as the lowest order of an expansion in the coupling strength [10]. This is not so surprising as a weak coupling implies weak forces and hence small accelerations.

3. Non-instantaneous and instantaneous action-at-a-distance

By identifying the following expectation values of field operators with their representation in terms of test particles

$$g_S^2 \int d^4x\, d^4y\, \left\langle \bar\psi(x)\psi(x)\,G_S(x-y)\,\bar\psi(y)\psi(y)\right\rangle \;\rightarrow$$
$$\hat{g}_s^2 \sum_{\substack{i,j=1\\ i\neq j}}^{A\cdot N} \int d\tau_i d\tau_j\; G_S(r(i,\tau_i) - r(j,\tau_j)) \tag{19}$$

$$g_V^2 \int d^4x\, d^4y\, \left\langle \bar\psi(x)\gamma^\alpha\psi(x)\,G_V(x-y)\,\bar\psi(y)\gamma_\alpha\psi(y)\right\rangle \;\rightarrow$$
$$\hat{g}_V^2 \sum_{\substack{i,j=1\\ i\neq j}}^{A\cdot N} \int d\tau_i d\tau_j\; G_V(r(i,\tau_i) - r(j,\tau_j))\,u_\alpha(i,\tau_i)u^\alpha(j,\tau_j) \tag{20}$$

in the action (8) one obtains the *non-instantaneous* action-at-a-distance

$$\int d^4x\, \langle\mathcal{L}(x)\rangle \;\rightarrow$$
$$\begin{aligned} \mathcal{A} \;=\; & -\frac{1}{2}\sum_{i=1}^{A\cdot N}\int d\tau_i\left(\widehat{M} - \lambda(i,\tau_i)\right)u(i,\tau_i)^2 \\ & + \frac{1}{2}\hat{g}_S^2 \sum_{\substack{i,j=1\\ i\neq j}}^{A\cdot N}\int d\tau_i d\tau_j\; G_S(r(i,\tau_i) - r(j,\tau_j)) \\ & - \frac{1}{2}\hat{g}_V^2 \sum_{\substack{i,j=1\\ i\neq j}}^{A\cdot N}\int d\tau_i d\tau_j\; G_V(r(i,\tau_i) - r(j,\tau_j))u_\alpha(i,\tau_i)u^\alpha(j,\tau_j)\,. \end{aligned} \tag{21}$$

Langrange multipliers $\lambda(i,\tau_i)$ have been introduced to ensure $u(i,\tau_i)^2 = 1$ and $\widehat{M} = M/N$ is the proportionate rest mass of the test particle.

If one now chooses a synchronizing space–like surface $S(x)$ (for example $S(x) = n_\alpha x^\alpha$ with n_α being a fixed time–like unit vector) one can expand each world line around this surface as a straight line, see eq. (12) and fig. 1. The synchronized proper times $\tau_i^s(t)$ are determined from the condition that the test particles are at time t at this surface:

$$S\Big(r(i,\tau_i^s(t))\Big) = t \tag{22}$$

In the spirit of the small acceleration assumption discussed in the previous section one can use the straight line expansion in the action (21) and perform the integration over τ_j. This results in an *instantaneous* action–at–a–distance where the Lorentz–boosted Yukawa fields appear again and there is only one time, the *scalar* synchronizing time t.

$$\begin{aligned}\mathcal{A} = \int dt \sum_i \frac{1}{\partial_\alpha S(r(i,t))u^\alpha(i,t)} \Bigg[& -\frac{1}{2}\Big(\hat{M} - \lambda(i,t)\Big)\, u(i,t)^2 \\ & + \frac{1}{2}\sum_{\substack{j\\ j\neq i}} \frac{\hat{g}_S^2}{4\pi} \frac{\exp\{-\mu_S\sqrt{-R(i,j,t)^2}\}}{\sqrt{-R(i,j,t)^2}} \\ & - \frac{1}{2}\sum_{\substack{j\\ j\neq i}} \frac{\hat{g}_V^2}{4\pi} \frac{\exp\{-\mu_V\sqrt{-R(i,j,t)^2}\}}{\sqrt{-R(i,j,t)^2}}\, u_\alpha(i,t)u^\alpha(j,t) \Bigg]\end{aligned} \tag{23}$$

The four–positions $r^\alpha(i,t) \equiv r^\alpha(i,\tau_i(t))$ and four–velocities $u^\alpha(i,t) \equiv u^\alpha(i,\tau_i(t))$ are to be taken at the same scalar synchronizing time t and

$$R^\alpha(i,j,t) = (r^\alpha(i,t) - r^\alpha(j,t)) - \Big[(r_\beta(i,t) - r_\beta(j,t))u^\beta(j,t)\Big]\, u^\alpha(j,t)\ . \tag{24}$$

The small acceleration approximation together with the introduction of a synchronizing hypersurface $S(x)$ leads to an equal time Lagrangian which is Lorentz–scalar and written in a manifestly covariant way.

Giving up explicit covariance and choosing $n = (1,0,0,0)$ in a special coordinate frame, the positions and velocities take the form

$$r(i,\tau_i(t)) = \Big(t\,,\, \vec{r}(i,t)\Big) \quad \text{and} \quad u(i,\tau_i(t)) = \frac{1}{\sqrt{1-\vec{v}^{\,2}(i,t)}}\Big(1\,,\, \vec{v}(i,t)\Big)\ . \tag{25}$$

With that a Lagrange function $\mathcal{L}(\vec{r}(i,t), \vec{v}(i,t))$ can be defined which depends only on the independent variables and one time. Even the Lagrange multipliers $\lambda(i,t)$ are not needed anymore if the variation is with respect to $\vec{v}(i,t)$ instead of all four $u^\alpha(i,t)$.

The advantage of the instantaneous lagrangian is that one can define easily the hamiltonian and the total momentum, which are then strictly conserved by the equations of motion. The disadvantage is that the equations of motion which result from the action (23) are much more complicated than those given in eq. (17) where the fields have been approximated.

We tried both sets of equations and did not see a significant difference for 640 testparticles describing the mean field of the collision ^{16}O $+^{16}$O at $E_{lab} = 1$ AGeV.

4. Non–instantaneous action–at–a–distance with retarded Green functions for nucleon–nucleon scattering

In the last years Bush and Nix [11] developed a method based on the Walecka model which describes the non–quantal but highly relativistic scattering of two nucleons. They take the finite size of the nucleons into account and use it to avoid runaway solutions. The resulting equations of motion are of integro–differential–type and rather involved. In their case radiation occurs and shows up as energy loss of the classically treated nucleons. In fig. 2 we show a comparison of the scattering angle as a function of impact parameter between their result and first order PRM or the small acceleration approximation. One should pay attention to the fact that for impact parameters b less than about 1 fm the density of the nucleon in the overlapping tail is more than about 15%, so that an effective mean–field theory like the Walecka model becomes questionable. Although the assumption (13) of small accelerations is not at all fulfilled as can be seen from fig. 3 the final results are not so different, especially for $b > 1$ fm the scattering angles are rather similar for the different lab–momenta. This shows that the approximation is very robust and does not lead to completely unphysical results even when the smallness parameter $\sqrt{-a^\mu a_\mu}/\mu_S = 40$.

5. Final remarks

Different from the Coulomb potential the Yukawa potential has a finite range given by the inverse of the meson mass. Massive fields stay close to their sources and cannot be radiated off so easily from an accelerating charge. This makes the action–at–a–distance formulation of Schwarzschild, Tetrode and Fokker well suited for effective field theories used in nuclear physics. Formally a field theory with moving charges as source terms and the boundary condition of vanishing in- and outgoing free fields is equivalent to an action–at–a–distance formulation in which the fields are eliminated. All the information about the fields is in proper–time integrals along the world lines inside the past and future light cone. A reasonable numerical treatment is possible only if not all the history (and future) of the world lines is needed but only sections which are close to the space–time point where the field is to be calculated. Particles which undergo only small accelerations and interact by finite range potentials fulfill this condition best.

In the following I list three cases where Predictive Relativistic Mechanics could be applied for massive fields in hadron physics.

1) High energy nucleon–nucleon collision

If the treatment of runaway solutions by taking into account the finite size of the nucleons as proposed by Bush and Nix [11] is sufficiently accurate one could take their results to test the two lowest orders of PRM. As shown in this contribution first order is already close to their results. There are however limitations to the physical application due to the following two conditions which contradict each other. First one can only look at high relative momenta (small de Broglie wavelength) in order to regard the nucleons as classical particles. Second, the effective field theory is not adequate at short length scales where QCD degrees of freedom are resolved.

2) Mean–field effects in relativistic heavy–ion collisions

In collisions of nuclei a reliable treatment of the nuclear mean field is over a wide energy range essential to interpret and analyse the experimental data. I believe that as a step away

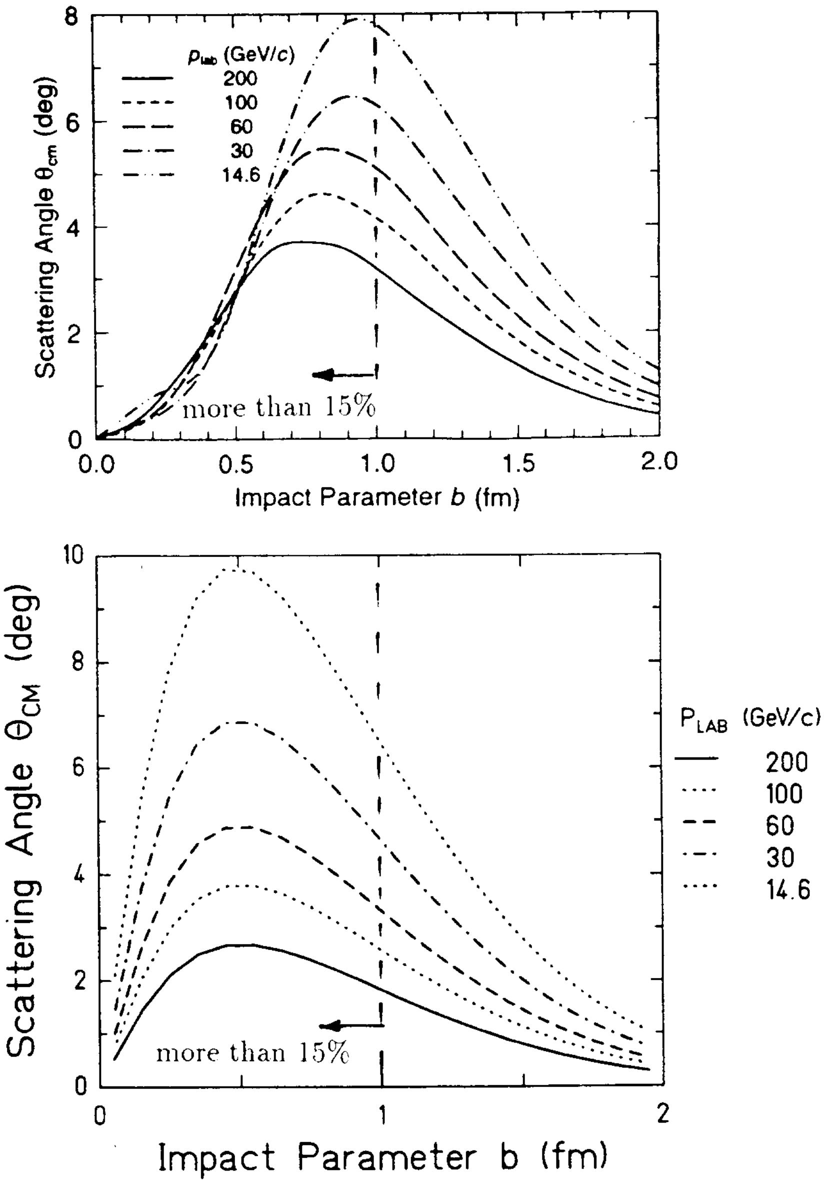

Figure 2
Scattering angle as function of impact parameter and lab momentum for nucleon–nucleon scattering. Upper part taken from Bush and Nix [11], lower part PRM first order approximation calculated by Jörg Lindner. For impact parameter below 1 fm the nucleon densities in the overlapping tails are more than 15% of their maximum.

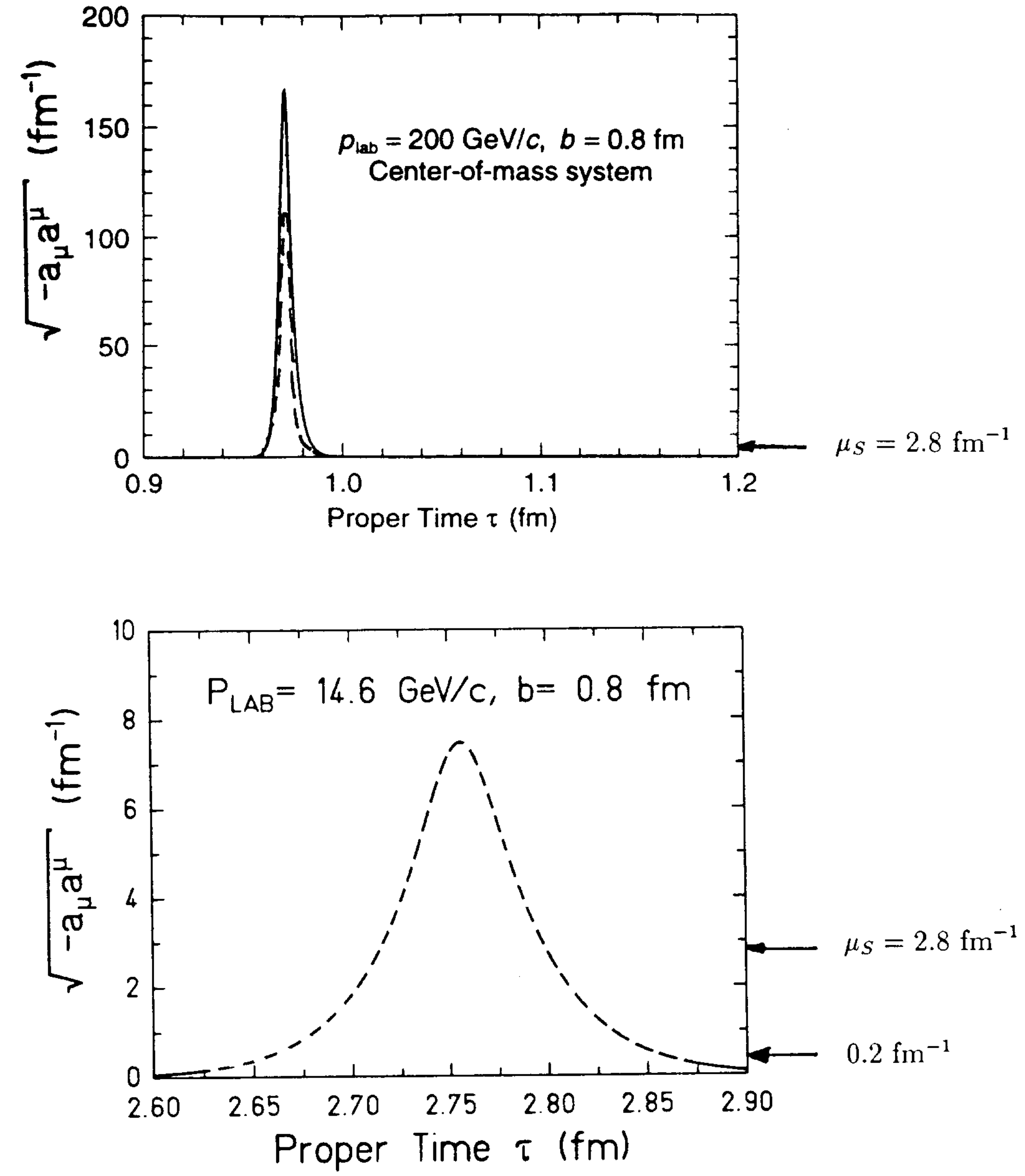

Figure 3
Acceleration of world lines for two different lab momenta, 200 GeV/c upper frame, 14.6 GeV/c lower frame. Dashed lines show PRM results calculated by Jörg Lindner, full line result of Bush and Nix [11]. Small acceleration assumption requires $\sqrt{-a_\mu a^\mu} \ll 2.8$ fm^{-1}. Typical accelerations due to mean field are 0.2 fm^{-1}.

from a local density approximation which does not include the finite range and retardation effects the application of PRM to test particles representing the phase–space density is most promising. However, due to the large number of test–particles only the lowest order can be treated numerically. First order is probably sufficient as the accelerations due to the mean field are not large. Questions still to be investigated are, first, can the Lorentz–boosted Yukawa potentials of eq. (14) and (16) in PRM in a consistent way be replaced by functions which do not diverge like $1/r$ for small distances r? This is necessary because only the mean–field part can be considered in an effective theory while hard collisions which occur below about 1 fm have to be treated differently. Second, can one define a total four–momentum, especially the energy, which is strictly conserved under the PRM equations of motion (and not only within the approximation) [12, 13]? Or can one modify the equations of motion without loosing their numerical simplicity such that a well defined total four–momentum is strictly conserved? This is highly desirable for checking the stability of numerical solutions. Third, is there an action–at–a–distance formulation when selfinteractions or field–dependent coupling strengths for the scalar field [14] are included?

3) Relativistic quantum mechanics

It seems that the long standing problem to calculate successfully saturation properties of nuclear matter from potentials which fit the free nucleon–nucleon phase shifts can be solved with the inclusion of relativistic effects into a G–matrix calculation [1]. Since long stationary nuclear systems have been studied [15] by eliminating the meson fields such that the hamiltonian has two–body terms like

$$\int d^3x \int d^4y \; \bar{\psi}(x)\psi(x)G(x-y)\bar{\psi}(y)\psi(y) \; , \tag{26}$$

where $\psi(x)$ is now the field operator for the nucleons. However, one is often using the static limit in which the integration over y^0 leads to an instantaneous Yukawa potential. Thus retardation effects are neglected.

Because of the success of the perturbative Feynman approach to QED the old motivation to formulate a Hamilton operator for interacting relativistic particles may be lost for electromagnetic fields but not for Yukawa fields. In hadron systems the situation is different in the sense, that pertubative treatments are usually not possible due to the strong coupling. A non–pertubative mean–field approach is the minimum requirement for a successful description [16] and because of the complex internal QCD structure of the hadrons very high precession cannot be expected anyhow.

In conclusion it might be worthwhile to reconsider action–at–a–distance formulations and methods of Predictive Relativistic Mechanics for solving effective field theories for hadrons.

I should like to express my thanks to Professor Luis Bel for very valuable discussions and comments.

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