An alternative to the gauge theoretic setting

Bert Schroer
CBPF, Rua Dr. Xavier Sigaud 150
22290-180 Rio de Janeiro, Brazil
and Institut fuer Theoretische Physik der FU Berlin, Germany

November 2010

Abstract

The standard formulation of gauge theories results from the Lagrangian (functional integral) quantization of classical gauge theories. A more intrinsic qunantum theoretical access in the spirit of Wigner's representation theory shows that there is a fundamental clash between the pointlike localization of zero mass (vector, tensor) potentials and the Hilbert space (positivity, unitarity) structure of QT. The quantization approach has no other way than to stay with pointlike localization and sacrifice the Hilbert space whereas the approach build on the intrinsic quantum concept of modular localization keeps the Hilbert space and trades the conflict creating pointlike generation with the tightest consistent localization:: semiinfinite spacelike string localization. Whereas these potentials in the presence of interactions stay quite close to associated pointlike field strength, the interacting matter fields to which they are coupled bear the brunt of the nonlocal aspect in that they are string.generated in a way which cannot be undone by any differentiation.

The new stringlike approach to gauge theory also revives the idea of a Schwinger-Higgs screening mechanism as a deeper and less metaphoric description of the Higgs spontaneous symmetry breaking and its accompanying tale about "God's particle" and its mass generation for all the other particles... .

1 Problematic aspects of the gauge theoretic setting of QED and QCD

The standard formulation of gauge theories follows the quantization parallelism to the classical Maxwell theory coupled to classical sources. A object of significant classical computational importance is the vectorpotential as it appears in the form of Lienard-Wiechert potentials of classical charge distributions. It behaves like a covariant classical vectorfield, apart from the fact that it cannot be uniquely fixed in terms of Cauchy data. Except for this particularity it behaves as a classical covariant causally propagating pointlike field.

Its quantized form entered the discourse of QFT almost from its beginnings, and by 1929 the status of gauge theories has been competently expressed in a review of the first phase of QFT [1][2].

As we know nowadays, the imperfections were not that insignificant as the author thought at that time, otherwise it would be difficult to understand that it took another two decades to use indefinite metric Krein spaces in the Gupta-Bleuler formalism[1] in order to make the quantum vectorpotentials compatible with the structure of state space in QT; the nonabelian gauge theories required a more complicated operator-ghost setting which was developed much later and which bears the name of the initials BRST of the protagonists Becchi, Rouet, Stora and Tuytin. These technical additions were important for the renormalization project and the return to physical observables, but they only affect some technical aspects of the quantization transfer of the classical gauge concept and not the idea as such. In fact apart from spacetime symmetries it is the only "symmetry" in classical physics; the "inner" symmetries (flavour symmetries) are of pure QFT origin and are a consequence of the superselection theory of localizable charges [3]. They are often red back into classical field theory with the purpose to present a QFT model as a canonical or functional integral quantization.

In both cases one was forced to do the renormalization calculations involving covariant pointlike fields in an indefinite metric setting; only at the end one could expect to return to a physical Hilbert space via a GNS construction based on the positive-semidefinite gauge invariant correlation functions. The construction of non pointlike matter fields carrying a Maxwell- or Yang-Mills- charge and their physical correlations, including those of matter fields, remained an extremely difficult if not impossible task and.

It is not without irony that the post QED formalism for strong interactions (the meson-nucleon interaction) of the 60s was conceptually much simpler. With its mass gaps it represented precisely the arena in which all the prerequisites for the derivation of the LSZ scattering theory including the Haag-Ruelle derivation from the locality and spectral principles in a Wightman setting hold; in fact it can hardly be called an accident that modern scattering theory in QFT was developed at the time of strong interactions between nucleons and mesons.

The attempts to include QED into this ideal world of relativistic particle physics (described by "interpolating" local quantum fields) failed[2]. The first difficulties were noted in form of perturbative infrared divergences in scattering theory; they were bypassed by constructing quantities which are apparently infrared convergent as the photon-inclusive cross sections for scattering of charged particles [4][5]. These calculations received also nonperturbative support from

[1] The application of the canonical quantization rules to models involving vector-fields was always a point of contention and, as will be shown here, the indefinite metric gauge theoretic formulation is by no means the last word.

[2] The occurrence of infrared divergences in a model of QFT is always an indication that there exist not yet understood spacetime localization properties. Even though they can be patched up by momentum space manipulations as the abandonment of scattering amplitudes in favor of certain inclusive cross sections, there is, in the long run, no substitute for a direct understanding in terms of spacetime localization properties.

the observation that the Källén Lehmann spectral functions of infraparticles which start with a cut singularity at $p^2 = m^2$ instead of a mass-shell delta function, in that case the LSZ limits vanish. This is precisely what one expects of the energy-momentum spectrum of indecomposable string-localized fields i.e. string-localized fields which cannot be written as an integral along a string. All these observations receive support from rigorous studies of the quantum Gauss law [3].

Similar to the YFS infrared calculation [5], these cut singularities which start at the mass shell are the result of summing of log singularities which leads to a coupling dependent power cut. Since spectral functions are positive measures, the strength of the cut must be smaller then that of the delta function; the large p behavior is however not limited beyond the temperedness restriction of the two-point distribution.

The reason why the construction of Gupta-Bleuler factor spaces, respectively in case of BRST the construction of the cohomology space (related to invariance under the nonlinear BRST transformation) leads to the unsolved problems of charged matter is that precisely in this step the non-compact localized physical charged operators have to emerge from their pointlike unphysical gauge dependent pointlike counterparts. Maxwell quantum charge and its particular type of semiinfinite stringlike generators are synonymous. In the indefinite metric description the localization of gauge variant operators had no physical meaning, it only offers the technical advantage of being able to use the renormalization machinery of pointlike covariant fields. In fact in using pointlike vectorpotentials one has to be extremely careful in order to avoid incorrect conclusions about the absence of the quantum field theoretical Aharonov-Bohm effect or the absence of Maxwell charges (see next section).

Instead of the quantization approach, which inevitably leads into the well-known gauge theoretic formulation with all its physical problems and technical attraction, we will set up an intrinsic framework which results from combining Wigner's representation theoretical setting with the concept of modular localization (next section), both having nothing in common with any classical quantization parallelism. This leads automatically to string-localized vectorpotentials acting in a Hilbert space with scale dimension $d_{scale} = 1$ which is the power counting prerequisite for the existence of renormalizable interactions. Different from the pointlike gauge theoretic formulation in the indefinite metric setting which is required by quantization, the stringlike potential setting has the great advantage of permitting an extension to interactions involving higher spin $s > 1$ zero mass fields: there always exist a tensor potential with $s_{scale} = 1$ and a coupling to matter fields which fulfills the power counting requirement at least for certain couplings to low spin matter and to its own copies (viz Yang-Mills models).

In this setting the notion of gauge has been replaced by the notion of "causal, (modular[3]) localization". The results of both approaches agree on local observ-

[3]"Modular localization" stands for the intrinsic (field-coordinatization-independent) formulation of causal relativistic localization in QT. A formulation of QFT which highlights this

ables (always pointlike generated), on that subalgebra there holds in fact the identity: local observables = gauge invariant subalgebra. But on string-like objects (as matter fields carrying Maxwell charges) which are genuinely noncompact string-localized, the gauge formulation offers an unphysical formalism without a practicable way to construct the important physical quantities which cannot be pointlike generated. With the exception of selfcouplings between tensorpotentials (Yang-Mills, selfcouplings between tensor potentials $g_{\mu\nu}(x,e)$ etc.), the string-localization of the potentials remains quite harmless[4]; but their mild nonlocal changes (moderated by the linear connection to pointlike field strengths), affects the localization of matter fields in a radical way which can not be undone by any linear operation. Whereas the return of stringlike potentials to pointlike fields is quite simple (linear) in QED and somewhat less simple (nonlinear) in Yang-Mills theories, the violent reprocessing of pointlike localized unphysical matter fields into noncompact string-localized charged fields is a change which the gauge approach is incapable to describe.

The only formal similarity between the two settings is that the gauge changing transformations resemble the changes of vectorpotentials under changes of string directions $e \to e'$ (see end of next section).

In the next section the Wigner theory will be extended in such a way that the full spectrum of covariant possibilities is realized. This is achieved by resolving the clash between localization and Hilbert space in favor of maintaining the latter. This will, with the help of modular localization theory, automatically result in semiinfinite string-localized vectorpotentials. Important explicit formulas in terms of string-localized intertwiners from the unique Wigner representation to covariant string potentials will be presented. It also will be shown that the use of vectorpotentials in Stokes theorem applied to magnetic surface fluxes gives the correct Aharonov-Bohm effect only in case of use of the stringlike physical vectorpotential.

The last section comments on interactions and indicates how the Higgs mechanism allows a more physical presentation in terms of the "Schwinger-Higgs screening" which is a QFT localization analog of the quantum mechanical Debeye screening in which the Coulomb potential passes to an effective short range Yukawa potential which in the localization analogy means a pointlike re-localization of string-localized charge generators via screening.

The reader who looks for a more detaild account of some of the issues is referred to [6].

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property is often called "local quantum physics" (LQP).

[4]There are linear differential operators which reconvert the potentials into pointlike field strengths. For the generation of an irreducible set of operators the field strength are sufficient, but for formulating the renormalizable interaction one needs the string-localized potentials.

2 String- localization in Wigner representations from modular localization

The representation-theoretical approach to covariance and localization has been described before [7][8][6], therefore the following presentation will be short and problem-oriented.

Being confused by the many ad hoc linear field equations for relativistic wave functions which appeared in the 30s, some of them being equivalent to each other in their physical content, Wigner in 1939 published a completely intrinsic description in terms of irreducible unitary ray representations of the Poincaré group. He found that there are 3 classes of positive energy representation one massive class with halfinteger spin and two massless classes, the finite helicity class and infinite helicity class (originally called "continuous spin" class); the only representation theoretical difference between the two is the representation theory of its noncompact little group $E(2)$, the euclidean group in two dimensions which only in the last case is faithfully represented (its "translations" are nontrivial).

For the first two Wigner and in more detail Weinberg [9], found that there are many covariant field equations whose covariant representations in the dotted/undotted $(A, \dot{B})$ representation formalism are connected to the physical spin/helicity s by

$$\left|A-\dot{B}\right| \leq s \leq \left|A+\dot{B}\right|, \; m>0 \tag{1}$$

$$s=\left|A-\dot{B}\right|, \; m=0 \tag{2}$$

where the second line expresses the fact that, as a result of the nonfaithful representation of the fix point group $E(2)$ of a lightlike vector, one looses a large number of covariant generators, including the possibility of generating the (m=0,s=1) representation by a covariant vectorpotential. The third infinite spin representation class is the most exotic one; as a result of a continuous euclidean mass-like parameter κ it is, similar to the massive family, rather big; its physical content was beginning to get unravelled very recently recently [8]. In the two finite spin cases one can calculate the generating covariant fields in terms of intertwiner $u^{A\dot{B}}(p)$ which intertwine between the Wigner representation and its various spinorial representations by group theoretical methods [9]; for the infinite spin case there are no pointlike free fields rather all free fields are string-localized and the intertwiners are infinite component objects $u(p,e)$ which depend on a string direction e.

For zero mass finite spin representations (2) the standard covariantization only leads to field strength as $F_{\mu\nu}(x)$ for s=1, $R_{\mu,\nu,\kappa,\lambda}$ for s=2, etc. as the pointlike fields with the lowest possible dimension $d(s)=s+1$. The missing A_{μ}-vector and $g_{\mu\nu}$-tensor potentials do not appear in the covariantization of the Wigner representation approach. This is a manifestation of a clash between localization, in this case pointlike localization, and the Hilbert space setting of QT. The gauge theoretical setting results from this clash by relaxing the Hilbert space structure and retaining the pointlike localization. In this case the $A_\mu(x)$ lives in an "ghostly" indefinite metric space (Gupta Bleuler, BRST) and the physical Wigner-Fock space of the field strength is only recovered by passing to a factor- or cohomology space. In this setting the vectormeson is a formal device which guaranties the continued use of the standard pointlike field formalism. The alternative way of resolving the Hilbert space – localization clash is to keep the Hilbert space structure and relax on pointlike localization. It turns out that the only localization which never leads to problems with the Hilbert space structure is semiinfinite stringlike i.e. on a halfline: $x+\mathbb{R}_+e$ with e being the spacelike string direction. Allowing this weaker localization one can construct covariant string-localized vectorpotentials which live in the same Hilbert space as the pointlike field strength, which are related by differential operators, and which act cyclically on the vacuum.

Their intertwiners $u(p,e)$ are not easily accessible by group theoretical methods, here the method of modular localization is more efficient. This method combines the Wigner representation theoretical setting with that of modular localization.

First some general remarks about localization. Quantum theory comes with two notions of localization, the Born-Newton-Wigner localization which corresponds to a classical *action at a distance* dynamics describing processes with unlimited basic propagation speed. Its quantum reign is QM[5] and it comes with a position operator $\vec{x}_{op}$ whose spectral decomposition leads to wave functions $\psi(\vec{x})$ and their Born probability density $|\psi(\vec{x})|^2$ Finite velocities as the sound velocity in case of a system of coupled oscillators are effective velocities showing up in expectation values in certain states; strictly speaking they become exact in the asymptotic scattering limit.

This probabilistic localization continues to be present in QFT, but another more fundamental localization takes over, the modular localization. One may view it as the quantum counterpart of the causal Cauchy propagation of classical hyperbolic differential equations. In contradistinction to the BNW localization which is consistent with QM but not intrinsic to it (Born introduced it several years after the discovery of QM) the modular localization is completely intrinsic. This explains why mathematicians discovered it independently in the form of the Tomita-Takesaki modular theory of operator algebras whereas physicists came to it through their study of thermal properties (KMS states) of open systems. The connection with localization came later and obtained a helping hand from the thermal aspects of the localization in front of black hole horizons. In interacting QFT there are no particles in compact spacetime regions, the inexorable presence of interaction caused vacuum polarization as an epiphenomenon of modular localization prevents the existence of such localized one particle operators and states whereas with the frame-dependent BNW localization there are no such properties. Nevertheless both localizations coalesce in the asymptotic

[5]This includes relativistic QM [10] in the form of *direct particle interactions* (DPI) which is only asymptotically frame-independent but leading to an invariant S-matrix.

scattering region, which is crucial for the existence of a vacuum-polarisation-free frame-independent S-matrix for particles.

Although historically the operator algebraic formulation of modular theory exited before the (quite recent) modular localization of states, it is easier and more convenient to present the latter before the former.

The simplest context for a presentation of the idea of modular localization is in the context of the Wigner representation theory of the Poincaré group. It has been realized by Brunetti, Guido and Longo [7] and in a more special context in [11] there is a natural localization structure on the Wigner representation space for any positive energy representation of the proper Poincaré group.

The starting point is an irreducible representation U_1 of the Poincaré group including the antiunitary TCP reflection on a Hilbert space H_1 that after "second quantization" becomes the single-particle subspace of the Hilbert space (Wigner-Fock-space) H_{WF} of the associated interaction free QFT[6]. The construction proceeds according to the following steps [7][12][8].

One first fixes a reference wedge region, e.g. $W_0 = \{x \in \mathbb{R}^d, x^{d-1} > |x^0|\}$ and considers the one-parametric L-boost group (the hyperbolic rotation by χ in the $x^{d-1} - x^0$ plane) which leaves W_0 invariant; one also needs the reflection j_{W_0} across the edge of the wedge which is apart from a π-rotation in the transverse plane identical to the TCP transformation. The Wigner representation U_1 in H_1 is then used to define two commuting wedge-affiliated operators

$$\mathfrak{d}_{1,W_0}^{it} = U_1(0, \Lambda_{W_0}(\chi = -2\pi t)),\ J_{1,W_0} = U_1(0, j_{W_0}) \tag{3}$$

where attention should be paid to the fact that in a positive energy representation any operator which inverts time is necessarily antilinear[7]. A unitary one-parametric strongly continuous subgroup as Δ_{1,W_0}^{it} acting on H_1 can be written in terms of a selfadjoint generator as $\Delta_{1,W_0}^{it} = e^{-itK_{W_0}}$ and therefore permits an "analytic continuation" in t to an unbounded densely defined positive operators Δ_{1,W_0}^{s} with dense domains which decrease with increasing s. Poincaré covariance allows to extend these definitions to wedges W in general position, and intersections of wedges lead to the definitions for general localization regions (see later). Since the localization is clear from the context, a generic notation without subscripts will cause no confusion. With the help of this operator one defines the unbounded antilinear operator S_1 which has the same dense domain.

$$S_1 = \Delta_1^{\frac{1}{2}} J_1,\ dom S_1 = dom \Delta_1^{\frac{1}{2}} \tag{4}$$

$$J_1 \Delta_1^{\frac{1}{2}} J_1 = \Delta_1^{-\frac{1}{2}} \tag{5}$$

Whereas the unitary operator Δ_1^{it} commutes with the reflection, the antiunitarity of the reflection J_1 causes a change of sign in the analytic continuation

[6]The construction works for arbitrary positive energy representations, not only irreducible ones.

[7]The wedge reflection j_{W_0} differs from the TCP operator only by a π-rotation around the W_0 axis.

as written in the second line. This leads to the involutivity of the S-operator as well as the identity of its range with its domain

$$S_1^2 \subset \mathbf{1}$$
$$dom\ S_1 = ran\ S_1$$

Idempotent means that the S-operator has ± 1 eigenspaces; since it is antilinear the +space multiplied with i changes the sign and becomes the - space; hence it suffices to introduce a notation for just one of the two eigenspaces

$$K_1(W) = \{domain\ of\ \Delta_{1,W}^{\frac{1}{2}},\ S_{1,W}\psi = \psi\} \tag{6}$$
$$J_{1W}K_1(W) = K(W') = K_1(W)',\ duality$$
$$\overline{K_1(W) + iK_1(W)} = H_1,\ K_1(W) \cap iK_1(W) = 0$$

It is important to be aware that, unlike QM, we are dealing here with real (closed) subspaces K of the complex one-particle Wigner representation space H_1.

An alternative which avoids the use of real subspaces is to directly work with complex dense subspaces as in the third line. Introducing the graph norm of the dense space, the complex subspace in the third line becomes a Hilbert space in its own right. The upper dash on regions in the second line denotes the causal disjoint (which is the opposite wedge) whereas the dash on real subspaces means the simplectic complement with respect to the simplectic form $Im(\cdot,\cdot)$ on H_1.

The two equations in the third line are the defining property of what is called the *standardness* of a subspace[8]; any standard K-space permits to define an abstract s-operator as follows

$$S_1(\psi + i\varphi) = \psi - i\varphi \tag{7}$$
$$S_1 = J_1\Delta_1^{\frac{1}{2}}$$

whose polar decomposition (written in the second line) returns the two modular objects Δ_1^{it} and J_1 which outside the context of the Poincaré group has in general no geometric significance. The domain of the Tomita S-operator is the same as the domain of $\Delta^{\frac{1}{2}}$ namely the real sum of the K space and its imaginary multiple. Note that in the present context this domain is determined solely by Wigner's group representation theory without any reference to a (nonexistent) covariant position operator or an extrinsic probability notion.

It is easy to obtain a *net of K-spaces* and associated S operators by $U_1(a,\Lambda)$-transforming the K-space for the distinguished W_0. A bit more tricky is the construction of sharper localized subspaces via intersections

$$K_1(\mathcal{O}) \equiv \bigcap_{W \supset \mathcal{O}} K_1(W) \tag{8}$$

[8]According to the Reeh-Schlieder theorem a local algebra $\mathcal{A}(\mathcal{O})$ in QFT is in standard position with respect to the vacuum i.e. it acts on the vacuum in a cyclic and separating manner. The spatial standardness, which follows directly from Wigner representation theory, is just the one-particle projection of the Reeh-Schlieder property.

where $\mathcal{O}$ denotes a causally complete smaller region (noncompact spacelike cone, compact double cone). Intersection may not be standard, in fact they may be zero in which case the theory allows localization in W (it always does) but not in $\mathcal{O}$. Such a theory is still causal, but has no pointlike localized generators not local in the sense that its associated free fields are pointlike. However spacelike cone intersections, whose core is a semiinfinite spacelike string $x+\mathbb{R}_{+}e$ are always standard [7], which implies that semiinfinite string generating wave functions (wave function-valued distributions in x and e) exist in every positive energy representation. For the first two classes the K_1-space is standard for arbitrarily small $\mathcal{O}$, but this is definitely not the case for the infinite helicity family for which the compact localization spaces turn out to be trivial[9].

It is well known that there are two equivalent ways to get from Wigner representation spaces to interaction-free operators acting in a Wigner-Fock Hilbert space. The most intrinsic one is the direct functorial relation of the real subspaces $K_1(\mathcal{O})$ to spacetime-indexed net of operator algebras $\mathcal{A}(\mathcal{O}),\ \mathcal{O}\subset\mathbb{R}^4$. This functorial relation is often misleadingly called $2nd$ quantization, ignoring that (Lagrangian) quantization is an art and not a mathematical functor. This functorial relation is as unique as the Wigner representation theory; the large choice of possible coordinatizations of these algebras in terms of operator-valued distributional covariant generators is related to the physical spin s by the formulae (1,2); but these are only the linear generators of the operator algebras, in addition there are infinitely many composite generators which can be expressed in terms of Wick-ordered polynomials of the linear ones. In view of the analogy with the coordinatization of geometry these operator-valued distribution generators constitute coordinatizations of spacetime-indexed nets of operator algebras.

The second way is better known [9] since it follows directly from the covariantization of Wigner wave functions. The result in the standard case is the well known formula for free fields in terms of the u,v intertwiners between the Wigner representation and their various covariant $(A,\dot{B})$ spinorial realizations

$$\Psi^{(A,\dot{B})}(x)=\frac{1}{(2\pi)^{\frac{3}{2}}}\int(e^{-ipx}\sum_{s_3=\pm s}u^{(A,\dot{B})}(p)\cdot a(p)+e^{ipx}\sum_{s_3=\pm s}v^{(A,\dot{B})}(p)\cdot b^{*}(p))\frac{d^3p}{2\omega} \tag{9}$$

Here a,b are the Wigner annihilation,creation operators which depend in addition to p on the little Hilbert space which is a representation space of the little group SU(2) or for m=0 E(2), the dot denotes the inner product in this Hilbert space. As mentioned $dim\Psi^{(A,\dot{B})}\geq s+1$,for $s\geq 1$ i.e. no higher spin matter has interactions within the boundaries of the power counting renormalizability.

There are explicit formulas for the intertwiners which have been derived by group theory, a method which is however not sufficient in case of string-localized covariant representations as needed for the (vector, tensor) potentials

[9] It is quite easy to prove the standardness for spacelike cone localization (leading to singular stringlike generating fields) just from the positive energy property which is shared by all three families [7].

in case of m=0. In that case the analog formulas have been derived in the setting of modular localization. The result can be written in a similar form. For $(m=0, s\geq 1)$ one finds for the covariant string-localized potentials with $\left|A+\dot{B}\right| \geq s \geq \left|A-\dot{B}\right|$ and the pointlike solution $s=\left|A-\dot{B}\right|$ in terms of field strengths excluded

$$\Psi^{(A,\dot{B})}(x;e) = \frac{1}{(2\pi)^{\frac{3}{2}}}\int(e^{-ipx}\sum_{s_3=\pm s}u^{(A,\dot{B})}(p,s_3;e)a(p,s_3)+ \tag{10}$$
$$+e^{ipx}\sum_{s_3=\pm s}v^{(A,\dot{B})}(p,s_3;e)b^*(p,s_3))\frac{d^3p}{2\omega}$$

with explicit formulas for the intertwiners which may be pictured as $(2A+1)(2B+1)\times 2\quad p,e$-dependent rectangular matrices which act on the two-component (labeled by helicities) columns of creation/annihilation operators. The only case which is important for the next section is the vectorpotential s=1

$$A^\mu(x,e) = \frac{1}{(2\pi)^{\frac{3}{2}}}\int(e^{-ipx}\sum_{s_3=\pm1}u^\mu(p,s_3;e)a(p,s_3)+e^{ipx}\sum_{s_3=\pm1}\overline{u^\mu(p,s_3;e)}a^*(p,s_3))\frac{d^3p}{2\omega}$$
$$u^\mu(p,s_3;e)_\pm = \frac{i}{pe+i\varepsilon}\{(\hat{e}_\mp(p)e)p^\mu-(pe)\hat{e}^\mu_\mp(p)\}$$

where the $\hat{e}_\pm(p)$ denotes the two p-dependent photon polarization vectors. The notation is also appropriate for the higher (integer) spins which also can be expressed in terms of higher tensor powers two polarization vectors and a higher power $(pe+i\varepsilon)^{-s}$ of the string-localization causing factor. Once one arrives at the formulas for the interwiners one can, without knowing anything about modular localization theory, check is covariance and locality properties

$$U(\Lambda)\Psi^{(A,\dot{B})}(x,e)U^*(\Lambda) = D^{(A,\dot{B})}(\Lambda^{-1})\Psi^{(A,\dot{B})}(\Lambda x,\Lambda e) \tag{11}$$
$$\left[\Psi^{(A,\dot{B})}(x,e),\Psi^{(A',\dot{B}')}(x',e'\right]_\pm = 0,\ x+\mathbb{R}_+e >< x'+\mathbb{R}_+e'$$

As expected, the scaling degree of the potential is $d_{sca}(A^\mu(x,e)) = 1$ i.e. better than that of the field strength. The resulting two-point function is of the form [13]

$$\langle A_\mu(x;e)A_v(x';e')\rangle = \int e^{-ip(x-x')}W_{\mu\nu}(p;e,e')\frac{d^3p}{2p_0},\ \curvearrowright [A_\mu(x;e)A_v(x';e')] \tag{12}$$
$$W_{\mu\nu}(p;e,e') = -g_{\mu\nu} - \frac{p_\mu p_\nu(e\cdot e')}{(p\cdot e-i\varepsilon)(p\cdot e'+i\varepsilon)} + \frac{p_\mu e_\nu}{(e\cdot p-i\varepsilon)} + \frac{p_\nu e'_\mu}{(e'\cdot p+i\varepsilon)}$$

The presence of the last 3 terms is crucial for the Hilbert space structure; without them one would fall back into the indefinite metric and negative probabilities.

Since free potentials and free field strength are always related by linear differential operators, it is not surprising that the two-point functions of the potentials can be written as an inverse power in pe times a tensorial expression in p and $e's$.

Instead of the gauge transformation there is now a rule for the change of the vectorpotential under $e \to e'$

$$A^\mu(x,e) \to A^\mu(x,e') + \partial^\mu \Phi(x;e,e') \tag{13}$$
$$\Phi(x,e,e') = \int e_\mu A^\mu(x+te',e)dt$$

This law for the change of strings continues to be valid in interacting theories in which the relation between string-localized potentials and physical (pointlocal) field strength remains linear, it however suffers interaction dependent modifications for Yang-Mills interactions which are in a one to one relation with the pointlike nonlinear field strength; in fact independence of e is synonymous with pointlike localization just as in the setting of gauge theory where gauge invariance is synonymous with pointlike generated local observables.

Already in the absence of interactions the unqualified use of the gauge formalism can lead to wrong results which are avoided in the Hilbert space string-localized setting. A famous case is the Aharonov-Bohm[10] effect or rather its specification in the setting of QFT. From the two-point function (12) one gets the commutator commutator of the stringlike potentials and from the latter one finds the commutator of the pointlike field strengths which are independent of e. We only need the equal-time restriction of the $H-E$ commutator

$$[H_i(\mathbf{x}), E_j(\mathbf{x}')] \sim \varepsilon_{ijk} \partial^k \delta(\mathbf{x}-\mathbf{x}') \tag{14}$$

from which one can compute the commutator of two ρ-regularized electric and magnetic delta function fluxes going through two orthogonal disks D_1 and D_2 which intersect in such a way that the boundary of one passes through the center of the other.

Following [14] one looks at a situation of two spatially separated, but interlocking regions $\mathcal{T}_1$ and $\mathcal{T}_2$ in which one represents as the smoothened boundary of two orthogonal unit discs. The delta function fluxes through the D_i are smoothened by convoluting $\star$ with a smooth function $\rho_i(\mathbf{x})$ supported in an ε-ball B_ε; the interlocking $\mathcal{T}_i$ are then simply obtained as $\mathcal{T}_i = \partial D_i + B_\varepsilon$ $i = 1,2$. One computes the following objects

$$[\vec{E}(\vec{g}_1)\vec{H}(\vec{g}_2)] = \int \vec{g}_1(x) rot \vec{g}_2(x) d^3x = \int \rho_1(x)d^3x \int \rho_2(y)d^3y \tag{15}$$
$$= \int \rho_1(x)d^3x \int \rho_2(y)d^3y \quad \vec{g}_i = \vec{\Phi}_i \star \rho,\ \vec{\Phi}_i(\vec{f}) = \int_{D_i} \vec{f} d\vec{D}_i$$

[10] The A-B effect is a semiclassical effect of quantum mechanical matter in an external magnetic potential.

The Φ_i is the functional which describe the flux through D_i, a kind of surface delta function. Of one would use the curl relation between the magnetic field and the unphysical pointlike vectorpotential the application of the Stokes theorem would lead to an integral over the $\mathcal{T}_i$ which vanishes. On the other hand the the same integral in terms of the string-localized potential gives the correct nonvanishing A-B expression.

This does not mean that the gauge theoretic setting is wrong, but only that one has to be careful. The safest strategy would be to pass after the calculations are finished and before the physical interpretation to the gauge invariant objects; but the difficulty in achieving this for the non pointlike generated charged matter is the motivation for writing this paper.

3 Stringlike potentials in interactions, Schwinger-Higgs charge screening

There are very different ways to introduce interactions; one completely intrinsic nonperturbative method which starts with classifying generators of wedge algebras whose properties are very closely related to crossing and analytic exchange properties of the bootstrap-formfactor program and works its way down to intersections by forming intersections which lead to arbitrarily small double cone algebras and their pointlike field generators. This construction method is very much at its beginnings and has only been understood in the case of factorizing models in d=1+1. In those cases where it works it leads to an existence proof and an explicit construction, an achievement which no other method has attained since the beginning of QFT.

Its intuitive basis is the insight that the weaker the spacetime localization, the better the control of the ubiquitous vacuum polarizations in the presence of interactions. Next to the whole spacetime, in which it is possible to find all operators, including those which applied to the vacuum create pure one-particle states (without accompanying vacuum polarization clouds), the best compromise between field (or operator algebra) localization and particles in the presence of (any kind of) interactions is the (noncompact) wedge region W. Of pivotal importance for this method is the relation between the modular objects and the scattering matrix ([6])

QFT's in d=1+3 can presently only be accessed by perturbing free fields with polynomial interaction densities. In this case the modular localization method can only be used for free fields. As explained in the previous section this only leads to new insights in case of $(m = 0, s \geq 1)$ representations i.e. in case of gauge theories and higher spin tensor potentials. As also explained there, the use of the string-localized potentials also solves the problem of keeping the free short distance dimensions at $d_{sd} = 1$ thus preventing their exclusion for reasons of nonexistence of renormalizable interactions. Unlike the previous nonperturbative method, perturbative series for correlation functions are known to always diverge i.e. they cannot be used to secure the mathematical existence of a model; at best they are asymptotically convergent for infinitesimally small couplings; properties which are true in every order are believed to indicate structural characteristics of the model.

Perturbation theory is usually formulated in terms of Lagrangian quantization or the closely related functional integral representations. But is has been known for a long time [15] that Feynman rules do not need free fields of the Euler-Lagrange type, any type of free field as they arise through covariantization from Wigner representations will do to be used in a scalar interaction density of the causal perturbation (Epstein-Glaser) setting. In fact the string-localized potentials of the previous section are definitely not Euler-Lagrange and can not be used in any quantization scheme be it Lagrangian or functional integral, hence the Euler-Lagrange propery is a restriction required by those special ways of accessing perturbation theory but not of perturbation theory per se.

In the previous section we learned that the full covariance spectrum (1) for zero mass finite helicity representation can be regained by admitting stringlike fields. The pointlike free *field strength* [11] is then connected with the free *stringlike potentials* by covariant differential operators. Both, the pointlike field strength and the stringlike potentials do not only create the same Hilbert space from the vacuum, they also fulfill the Reeh-Schlieder theorem (popular name: state-operator relation) which in case of string-localization means that the operators with a localization around an arbitrary small neighborhood of the (x, e) string applied to the vacuum is dense in the Hilbert space.

We have presented structural arguments in favor of using stringlike potentials (rather than pointlike field strength) even in the absence of interactions when Stokes argument is invoked for surface integrals over magnetic fluxes as in the QFT Aharonov-Bohm argument. The A-B effect is correctly described with string-localized potentials whereas pointlike potentials lead to a zero effect. It is also well-known that the Maxwell charges vanish in the standard gauge indefinite metric setting pointlike potentials and that the presence of indefinite metric prevents the validity of the Dirac-Maxwell equations [14][**6**]. This is another call to be careful with drawing physical conclusion in the gauge setting. The rule of gauge theory is of course to go first to the gauge-invariant correlations and perform a GNS reconstruction of the corresponding operators in the canonically associated Hilbert space. Only after having achieved this one can draw physical conclusions.

There is no problem with the subalgebra of pointlike (strictly observable) generators. But the Maxwell- or Yang-Mills- charge carrying operators are never pointlike generated and the gauge setting offers no strategy to construct them. These problems become particularly pressing if one looks at article and textbooks [16] on QCD where the technical advantage of the analytic continuation method of dimensional regularization with respect to gauge theories is misunderstood as making renormalized pointlike quark fields objects of physical interests.

[11] We use this terminology in a generalized sense; all the pointlike generators (the only ones considered in [9]) are called field strength (generalizing the $F_{\mu\nu}$) whereas the remaining string-localized generators are named potentials.

Of course the presence of the ghost degrees of freedom renders the gauge setting renormalizable. However the string-localization in a Hilbert space does the same: for each m=0 spin $s \geq 1$ there exists always a potential of lowest possible dimension namely $d_{sca}(\Psi^{(\frac{s}{2},\frac{\dot{s}}{2})}(x,e)) = 1$ which is the power-counting prerequisite for constructing renormalizable interactions. Strictly speaking the renormalizabilty of pointlike would already stop after $s = 0, \frac{1}{2}$; beyond there is only the alternative of either using the gauge approach (which is only known for m=0,s=1) or the string-localized potential setting for which there are renormalizable candidates for any s. The short distance dimension of pointlike objects increase with s; this is well known for the massive case where the covariance for pointlike fields covers the whole spinorial spectrum (1) and there is no need for string-localization coming from representation theory. The simplest example would be a massive pointlike vector field $A_\mu(x)$ with $d_{sd} = 2$ whereas the dimension of $A_\mu(x,\mu)$ is $d_{sd} = 1$. It is only the stringlike massive potential which has a massless limit.

It was already mentioned that the string-localization has hardly any physical consequences for photons, since even in the presence of interactions the content of the calculated theory can be fully described in terms of linearly related pointlike field strengths. Even the scattering theory of photons in the charge zero sector has no infrared problems and follows a similar logic as LSZ [17]. However the interaction-induced string-localization of the charged field which is transferred from the vectorpotentials[12] is a much more serious matter; it is inexorably connected with the electric charge, and there is no linear operation nor any other manipulation which turns the noncompact localization of charged quantum matter into compact localization; electrically charged operators have no better generators than string-localized ones. The argument [3] based on the use of the quantum adaptation of Gauss's law shows that the noncompact (at best stringlike) localization nature of generating Maxwell charge-carrying fields is not limited to perturbation theory. The stringlike localization is so strong that even the Lorentz symmetry becomes spontaneously broken in nontrivial charge sectors.

It is customary to refer to Maxwell or Yang-Mills theories as *local* gauge theories and to theories involving complex fields (scalar or spinor) and which have interactions which are invariant under constant phase or in case of multiplets under SU(N) groups (e.g. old-fashioned meson-nucleon models) as invariant under global gauge group, But this is a somewhat treacherous terminology which only refers to superficial formal aspects but ignores the deeper physical distinction. From the viewpoint of localization it is just the other way around, namely the superselection theory which leads to standard inner symmetries is built on compactly localizable charges (the DHR superselection theory [3]) whereas for noncompact string-localized charges a superselection theory, if it exists at all, has continuously many superselected sectors and its inner symmetry is unknown. In other words the beautiful reconstruction of superselection sectors and charge-

[12] Localization of the free fields, in terms of which the interaction is defined in the perturbative setting, is not individually preserved in the presence of interactions; the would be charged fields are not immune against delocalization from interactions with stringlike vectorpotentials.

carrying field algebras from their observable shadow (Marc Kac: how to hear the shape of a drum) is presently limited to compactly localizable (pointlike generated) charges.

Apart from photon field strength, fields are not directly observable; nobody has ever measured a hadronic field.

Its most dramatic observable manifestation occurs in the scattering of charged particles. As mentioned before, the infrared peculiarities of scattering of electrically charged particles, first noted by Bloch and Nordsiek, were observed at the time as the stringlike Dirac-Jordan-Mandelstam formula from gauge theory ([6]), but the two observations remained disconnected [13]. the standard perturbative gauge formalism (which existed in its non-covariant unrenormalized form since the time of the B-N paper) was not capable to address the construction of string-localized physical fields. This is particularly evident in renormalized perturbation theory which initially seemed to require just an adaptation of scattering theory [5], but whose long term consequences, namely a radical change of one-particle states ("infraparticles") and the ensuing loss of a tensor-factorizing Wigner-Fock as well as the spontaneous breaking of Lorentz invariance and a missing spin-statistics theorem for infraparticles, were much more dramatic.

Up to now the dramatic conceptual change was patched up by acting as if the theory is under the LSZ umbrella and counteracting the resistance of the theory against invalid assumptions by manipulating its resistance with infrared cutoffs (the spirit of "effective" QFT) and looking for infrared stable quantities which finally allow a removal of the cutoff (the photon inclusive cross section). But the last step, namely the direct connection of the inclusive cross section for charged infraparticles with an asymptotic limit of string-localized spacetime objects But the history of particle physics shows that whereas it is helpful to think up intelligent placeholders for conceptual problems which one cannot solve for the time-being, it is detrimental for progress to leave them up to the cows come home.

The spacetime setting in a theory as QFT, for which everything must be reduced to its localization principles, is much more important than in QM where stationary momentum space scattering formulations compete with time-dependent ones. As mentioned before Coulomb scattering in QM can be incorporated into any formulation of scattering theory by extracting a diverging phase factor which results from the long range. Noncompact string-localization is a more violent change from pointlike generated QFT than long- versus short range quantum mechanical interactions.

Perturbative scattering (on-shell) processes represented by graphs which do not contain inner photon lines turn out to be independent of the string direction e i.e. they appear as if they would come from a pointlike interaction[14]. This

[13]The DJM formula is outside the perturbative approach and does not reveal any physical reason why in a pointlike theory (Maxwell, Yang-Mills) charge has no pointlike generators, the other reason is that whereas LSZ scattering theory identifies momentum space scattering amplitudes directly with spacetime limits, the same procedure applied to charges fields gives zero and a direct spacetime procedure for inclusive cross sections is not known.

[14]The time-ordered correlation functions, of which they are the on-shell restriction, are

includes the lowest order Møller- and Bhaba scattering. The mechanism consists in the application of the momentum space field equation to the u, v spinor wave functions so that from (12) only the $g_{\mu\nu}$ term in the photon propagator survives. The terms involving photon lines attached to external charge lines do however depend on the string directions; these are the same graphs which in the old infrared investigations were responsible for the on-shell infrared divergences i.e. e-dependence and the graphical positioning of infrared behavior is synonymous; the on-shell infrared divergence and the distributional e-dependence of correlation functions which prohibits to put to e's on top of each other (short distance limit in the de Sitter space of the spacelike e-directions) the are two sides of the same coin. In fact the smearing in e and the careful constructions of "de Sitter composites" are the additional handles which the string formalism offers. This should give an interesting powerful separation of ultraviolet and infrared behavior which in particular in the standard gauge approach to QCD models has been a stumbling block. None of the gauge theoretic formalisms, not even the much celebrated dimensional regularization is capable to achieve this; the separation is only possible in the new string-localized setting.

In the sequel some remarks on the perturbative use of stringlike vectorpotentials for scalar QED are presented which is formally defined in terms of the interaction density[15]

$$g\varphi(x)^*(\partial_\mu\varphi(x))A^\mu(x,e) - g(\partial_\mu\varphi(x)^*)\varphi(x)A^\mu(x,e) \tag{16}$$

It is also the simplest interaction which permits to explain the Higgs mechanism as a QED charge-screening. The use of string-localized vectorpotentials as compared to the standard gauge formalism deflects the formal problems of extracting quantum data from an unphysical indefinite metric setting to the ambitious problem of extending perturbation theory to the realm of string-localized fields. This is not the place to enter a presentation of (yet incomplete) results of a string-extended Epstein-Glaser approach. Fortunately this is not necessary if one only wants to raise awareness about some differences to the standard gauge approach.

It has been known for a long time that the lowest nontrivial order for the Kallen-Lehmann spectral function can be calculated without the full renormalization technology of defining time-ordered functions. With the field equation

$$(\partial^\mu\partial_\mu + m^2)\varphi(x) = gA_\mu(x,e)\partial^\mu\varphi(x) \tag{17}$$

the two-point function of the right hand side in lowest order is of the form of a product of two Wightman-functions namely the point-localized $\langle\varphi(x)\varphi^*(y)\rangle = i\Delta^{(+)}(x-y)$ and that of the string-localized vectorpotential (12)

$$\langle A^\mu(x,e)A^\nu(x',e')\rangle \, \langle\partial_\mu\varphi(x)\partial_\nu\varphi^*(x')\rangle \tag{18}$$

however string-dependent.

[15]The integral over the interaction density is formally e-independent.

leading to the two-point function in lowest (second) order

$$(\partial_x^2+m^2)(\partial_{x'}^2+m^2)\left\langle \varphi(x)\varphi^*(x')\right\rangle_{e,e'}^{(2)} \sim g^2 \left\langle A^\mu(x,e)A^\nu(x',e')\right\rangle \left\langle \partial_\mu\varphi(x)\partial_\nu\varphi^*(x')\right\rangle \tag{19}$$

which is manifestly e-dependent in a way which cannot be removed by linear operations as in passing from potentials to field strength. One can simplify the e dependence by choosing collinear strings $e=e'$, but the vectorpotential propagator develops an infrared singularity and in general such coincidence limits (composites in d=2+1 de Sitter space) have to be handled with care (although these objects are always distributions in the string direction i.e. can be smeared with localizing testfunctions in de Sitter space); just as the problem of defining interacting composites of pointlike fields through coincidence limits. The infrared divergence can be studied in momentum space; a more precise method uses the mathematics of wave front sets[16]. This simple perturbative argument works for the second order two-point function, the higher orders cannot be expressed in terms of products of Wightman function but require time ordering and the Epstein-Glaser iteration.

Not all functions of the matter field φ are e-dependent; charge neutral composites, as e.g. normal products $N(\varphi\varphi^*)(x)$ or the charge density are e-independent. On a formal level this can be seen from the graphical representation since a change of the string direction $e \rightarrow e'$ (13) corresponds to an abelian gauge transformation. The divergence form of the change of localization directions together with the current-vectorpotential form of the interaction reduces the e-dependence of graphs to *vectorpotentials propagators attached to external charged* lines while all e-dependence in loops cancels by partial integration and current conservation. This is in complete analogy to the standard statement that the violation of gauge invariance and the cause of on-shell infrared divergencies on charged lines result from precisely those external charge graphs; external string-localized vectorpotential lines cause no problems since they loose their e-dependence upon differentiation. A *neutral external composite* as $\varphi\varphi^*$ on the other hand does not generate an external charge line; again the gauge invariance argument parallels the statement that such an external vertex does not contribute to the string-localization.

Hence both the gauge invariance in the pointlike indefinite metric formulation and the e-independence in the string-like potential formulation both lead to pointlike localized subtheories[17]. But whereas the embedding theory (Gupta-Bleuler, BRST) in the first case is unphysical, the string-like approach uses Hilbert space formulations throughout. The pointlike localization in an indefinite metric description is a fake. Its technical advantage is that pointlike interactions, whether in Hilbert space or in a indefinite metric setting, are treatable

[16] Technical details as renormalization, which are necessary to explore these unexplored regions, will be deferred to seperate work.

[17] Note however that the spacetime interpretation of the e is not imposed. The proponents of the axial gauge could have seen in in the free two-pointfunction of vectorpotentials and in all charge correlators if they would have looked at the commutators inside their perturbative correlation functions. The axial "gauge" is not a gauge in the usual understanding of this terminology.

with the same well known formalism. The gauge invariant correlation define (via the GNS construction) a new Hilbert space which coalesces with the subspace obtained by application of the pointlike generated subalgebra of the physical string-like formulation to the vacuum.

But whereas the noncompact localized charge-carrying fields are objects of a physical theory, it has not been possible to construct physical charged operators through Gupta-Bleuler formalism or BRST cohomological descent. The difficulty here is that one has to construct non-local invariants under the nonlinear but formally local acting BRST symmetry. So the simplicity of the gauge formalism has a high prize when it comes to the construction of genuinely nonlocal objects as charged fields.

This leaves the globally charge neutral *bilocals* in the visor. Their description is expected to be given in terms of formal *bilocals which have a stringlike "gauge bridge"* linking the end points of the formal bilocals. In contrast to the string-localized single operators it is difficult to construct them in perturbation theory starting from string-localized free fields; they are not part of the interacting free fields and they are too far removed from the form of the interaction. In order to understand the relation between such neutral bilocals and infraparticles one should notice that in order to approximate a scattering situations, the "gauge-bridge" bilocals will have to be taken to the limiting situation of an infinite separation distance, so that the problem of the infinite stringlike localization cannot be avoided, since it returns in the scattering situation. The only new aspect of the proposed approach based on string-localized potentials which requires attention is that the dependence on the individual string directions e is distributional i.e. must be controlled by (de Sitter) test function smearing and moreover that composite limits for coalescing $e's$ can be defined. But a separation of ultraviolet behavior in the Minkowski space x from the infrared aspects encoded in to the ultraviolet aspects of the de Sitter $e's$ is just what was missing in the standard approach.

Finally there is the problem of Schwinger-Higgs mechanism in terms of string localization. The standard recipe starts from scalar QED which has 3 parameters (mass of charged field, electromagnetic coupling and quadrilinear selfcoupling required by renormalization theory). The QED model is then modified by Schwinger-Higgs screening in such a way that the Maxwell structure remains and the total number of degrees of freedom are preserved. The standard way to do this is to introduce an additional parameter via the vacuumexpectation value of the alias charged field and allow only manipulations which do not alter the degrees of freedom. We follow Steinmann [18], who finds that the screened version consists of a selfcoupled real field R of mass M coupled to a vectormeson A^μ of mass m with the following interaction

$$L_{int} = gmA^\mu A_\mu R - \frac{gM^2}{2m}R^3 + \frac{1}{2}g^2 A_\mu A^\mu R^2 - \frac{g^2M^2}{8m^2}R^4 \tag{20}$$

$$\Psi = R + \frac{g}{2m}R^2 \tag{21}$$

The formula in the second line is obtained by applying the prescription $\varphi \rightarrow$

$\langle\varphi\rangle + R + iI$ to the complex field within the neutral (and therefore point-local) composite $\varphi\varphi^*$ and subsequently formally eliminating the I field by a gauge transformation. The result is the above interaction where A_μ and R are now massive fields. Since the field Ψ is the image of a pointlike $\varphi\varphi^*$ under the Higgs prescription, the real matter field Ψ, as the screened version of the pointlocal charge-neutral $\varphi\varphi^*$remains local. However the screening does not only break the charge symmetry (thus trivializing the charge) and uses the I degrees of freedom to convert the photon into a massive vectormeson, but also disrupts the even-odd symmetry $R \to -R$ of the remaining R-interaction. It is the absence of this $\mathbb{Z}_2$ selection rule which transfers the pointlike localization of Ψ to R so that together with the pointlike $F_{\mu\nu}$[18] from the stringlike $A_\mu(x,e)$ the screened model is generated in terms of only pointlike fields.

Hence in the present context the string-localized potentials, as well as the gauge theoretic BRST formalism, behave as a "catalyzer" which makes a theory amenable to renormalization. The former have the additional advantage over in the latter that the Hilbert space is present throughout the calculation.

One has to be careful in order not to confuse computational recipes with physical concepts. Nonvanishing vacuum expectations (one-point functions) are part of a recipe and *should not be directly physically interpreted*, rather one should look at the intrinsic observable consequences[19] before doing the physical mooring. The same vacuum expectation trick applied to the Goldstone model of spontaneous symmetry breaking has totally different consequences from its application in the Higgs-Kibble (Brout-Englert, Guralnik-Hagen) symmetry breaking.

In the case of spontaneous symmetry breaking (Goldstone), the *charge associated with the conserved current diverges* as a result of the presence of a zero mass Boson which couples to this current. On the other hand in the Schwinger-Higgs screening situation *the charge of the conserved current vanishes* (i.e. is completely screened) and hence there are no charged objects which would have to obey a charge symmetry with the result that the lack of charge resulting from a screened Maxwell charge looks like a symmetry breaking.

$$Q_{R,\Delta R} = \int d^3x j_0(x) f_{R,\Delta R}(x), \quad f_{R,\Delta R}(x) = \begin{matrix} 1 \; for \;\; |\mathbf{x}| < R \\ 0 \;\; for \;\; |\mathbf{x}| \geq R + \Delta R \end{matrix}$$

$$\lim_{R\to\infty} Q_{R,\Delta R}^{spon} |0\rangle = \infty, \; m_{Goldst} = 0; \quad \lim_{R\to\infty} Q_{R,\Delta R}^{screen} \psi = 0, \; all \; m > 0$$

That the recipe for both uses a shift in field space by a constant does not mean that the physical content is related. The result of screening is the vanishing of a Maxwell charge which (as a result of the charge superselection) allows a copious production of the remaining R-matter.

Successful recipes are often placeholders for problems whose better understanding needs additional conceptual considerations. In both cases one can

[18] From the pointlike $F_{\mu\nu}$ on can construct a pointlike $A_\mu(x)$ with the same dimension.

[19] These are properties which can be recovered from the observables of the model i.e. they do not depend on the particular prescription use in its construction.

easily see that the incriminated one-point vacuum expectation has no intrinsic physical meaning, i.e. there is nothing in the intrinsic properties of the observables of the two theories which reveals that a nonvanishing one-point function was used in the recipe for its construction. For a detailed discussion of these issues see [19].

The premature interpretation in terms of objects which appear in calculational recipes tends to lead to mystifications in particle theory; in the present context the screened charged particle has been called the "God particle". As mentioned before the Schwinger-Higgs screening is analog to the quantum mechanical Debeye screening in which the elementary Coulomb interaction passes to the screened large distance effective interaction which has the form of a short range Yukawa potential. The Schwinger-Higgs screening does not work (against the original idea of Schwinger) directly with spinor- instead of scalar matter. If one enriches the above model by starting from QED which contains in addition to the charged scalar fields also charged Dirac spinors then the screening mechanism takes place as above via the scalar field which leads to a loss (screening, bleaching) of the Maxwell charge while the usual charge superselection property of complex Dirac fields remains unaffected.

The Schwinger-Higgs mechanism has also a scalar field multiplet generalization to Yang-Mills models; in this case the resulting multicomponent pointlike localized massive model is much easier to comprehend than its "charged" string-localized origin. As the result of screening there is no unsolved confinement/invisibility problem resulting from nonabelian string-localization.

The Schwinger Higgs screening suggests an important general idea about renormalizable interactions involving massive $s \geq 1$ fields, namely that formal power-counting renormalizability ($d_{sd} = 1$) is not enough. For example a pure Yang-Mills interaction with massive gluons (without an accompanying massive real scalar multiplet) could be an incorrect idea because the string-localization of the Hilbert space compatible gluons could spread all over spacetime or there may exist other reasons why the suspicion that such theories are not viable may be correct. Such a situation would than be taken as an indication that a higher spin massive theory would always need associated lower spin massive particles in order to be localizable; in the s=1 case this would be the s=0 particle resulting from Schwinger-Higgs screening. Before one tries to understand such a structural mechanism which requires the presence of localizing lower spin particles it would be interesting to see whether these new ideas allow any renormalizable s=$\frac{3}{2}$ (Rarita-Schwinger) theories. Even though there may be many formal power-counting renormalizable massive $s \geq 1$ interactions only a few are expected to be pointlike localized.

It is interesting to mention some mathematical theorems which support the connection between localization and mass spectrum. The support for placing more emphasis on localization in trying to conquer the unknown corners of the standard model comes also from mathematical physics. According to Swieca's theorem [20][19] one expects that the screened realization of the Maxwellian structure is local i.e. the process of screening is one of reverting from the electromagnetic string-localization back to point locality together with passing from a gap-less situation to one with a mass gap. Last not least the charge screening leads to a Maxwell current with a vanishing charge[20] and the ensuing copious production of alias charged particles. The loss of the charge superselection rule in the above formulas (20) is quite extreme, in fact even the $R \leftrightarrow -R$ selection rule has been broken (20) in the above Schwinger-Higgs screening phase associated with scalar QED. The general idea for constructing renormalizable couplings of massive higher spin potentials interacting with themselves or with normal s=0,1/2 matter cannot rely on a Schwinger-Higgs screening picture because without having a pointlike charge neutral subalgebra for zero mass potentials as in QED, which is the starting point of gauge theory, there is no screening metaphor which could preselect those couplings which have a chance of leading to a fully pointlike localized theory, even though renormalizability demands to treat all $s \geq 1$ as stringlike objects with d_{sca}=1. Of course at the end of the day one has to be able to find the renormalizabe models which maintain locality of observables either in the zero mass setting as (charge-neutral) subalgebras (QED,Yang-Mills) or the massive theories obtained from the former with the help of the screening idea. gauge theory is a crutch whose magic power is limited to s=1, for s> 1 it lost its power and one has to approach the localization problem directly.

The existence of a gauge theory counterpart, namely the generalization of the BRST indefinite metric formalism to higher spins, is unknown. So it seems that with higher spin one is running out of tricks, hence one cannot avoid confront the localization problem of separating theories involving string-localized potentials which have pointlike generated subalgebras from those which are totally nonlocal and therefore unphysical. This opens a new chapter in renormalization theory and its presentation would, even with more results than are presently available, go much beyond what was intended under the modest title of this paper.

An understanding of the Schwinger-Higgs screening prescription in terms of localization properties should also eliminate a very unpleasant previously mentioned problem which forces one to pass in a nonrigorous way between the renormalizable gauge (were the perturbative computations take place) and the "unitary gauge" which is used for the physical interpretation. The relation between the two remains somewhat metaphoric.

The screened interaction between a string-localized massive vectorpotential and a real field (20) remains pointlike because the string localization of the massive vectorpotential only serves to get below the power counting limit but does not de-localize the real matter field; since the pointlike field strength together with the real scalar field generate the theory, the local generating property holds. In an approach based on string-localization there is only one description which achieves its renormalizability by string-localized potentials.

The BRST technology is highly developed, as a glance into the present literature [22] shows. It certainly has its merits to work with a renormalization formalism which starts directly with massive vectormesons [21] instead

[20]Swieca does not directly argue in terms of localization but rather uses the closely related analyticity properties of formfactors.

of the metaphoric "photon fattened on the Higgs one point function". It is hard to think how the BRST technology for the presentation of the Schwinger-Higgs screening model which starts with a massive vectormeson in [22] can be improved. For appreciating this work it is however not necessary to elevate "quantum gauge symmetry" (which is used as a technical trick to make the Schwinger-Higgs mechanism compliant with renormalizability of massive s=1 fields) from a useful technical tool to the level of a new principle.

Besides the Schwinger-Higgs screening mechanism which leads to renormalizable interactions of massive vectormesons with low spin matter, there is also the possibility of renormalizable "massive QED" which in the old days [23] was treated within a (indefinite metric) gauge setting in order to lower the short distance dimension of a massive vectormeson from $d_{sd} = 2$ to 1, and in this way stay below the power-counting limit. Such a construction only works in the abelian case; for nonabelian interactions the only way to describe interacting massive vectormesons coupled to other massive s=0,1/2 quantum matter is via Higgs scalars in their Schwinger-Higgs screening role. Whereas the local Maxwell charge is screened, the global charges of the non-Higgs complex matter fields are preserved. It seems that Schwinger's original idea of a screened phase of spinor QED cannot be realized, at least not outside the two-dimensional Schwinger model (two-dimensional massless QED).

But the educated conjectures in this section should not create the impression that the role of the Schwinger-Higgs screening in the renormalizability of interactions involving selfcoupled massive vectormesons has been completely clarified; if anything positive has been achieved, it is the de-mystification of the metaphor of a spontaneous symmetry breaking through the vacuumexpectation of a complex gauge dependent field and the tale of "God's particle" which creates the masses of s=1/2 quantum matter. Actually part of this de-mystification has already been achieved in [22].

This leads to the interesting question whether, apart from the presence of the Higgs particle (the real field as the remnant of the Schwinger-Higgs screening), there could be an intrinsic difference in the structure of the vectormeson. Such a difference could come from the fact that the screening mechanism does not destroy the algebraic structure of the Maxwell equation, whereas an interaction involving a massive vectormeson coming in the indicated way from a S-H screening mechanism and interacting with spinorial matter fields maintains the Maxwell structure. In the nonabelian case this problem does not arise since apparently the Schwinger-Higgs screening mechanism is the only way to reconcile renormalizability with localizability (or a return to physics from an indefinite metric setting).

This raises the interesting question whether renormalizability and pointlike locality of interactions with massive higher spin $s > 1$ potentials is always related to an associated zero mass problem via an analog of a screening mechanism in which a lower spin field plays the analog of the Higgs field?

Whereas for interactions between spin one and lower spin fields the physical mechanism behind the delocalization of matter (or rather its noncompact re-localization) is to some degree understood, this is not the case for interacting higher spin matter. Stringlike interactions enlarge the chance of potentially renormalizable (passing the power counting test) theories, in fact stringlike potentials with dimension $d_{sca} = 1$ exist for any spin (hence infinitely many) whereas the borderline for pointlike interaction is $s = 1/2$ and with the help of the gauge setting $s = 1$. Certain interactions, as the Einstein-Hilbert equation of classical gravity probably remain outside the power-counting limit even in the stringlike potential setting, but certain polynomial selfinteractions between the $g_{\mu\nu}(x, e)$ with $dim g_{\mu\nu}(x, e) = 1$ may be renormalizable. The existence of free pointlike field strength (in this case the linearized Riemann tensor) indicates that there may be renormalizable interactions which lead to pointlike subalgebras, but the presence of self-couplings modifies the transformation law under a change of e (13) which now depends on the interaction as it is well-known from the gauge theoretical formulation for Yang-Mills couplings.

One of course does not know whether QFT is capable to describe quantum gravity (it never has been tried), but if it does in a manner which is compatible with renormalized perturbation theory, there will be no way to avoid string-localized tensorpotentials even if the theory contains linear or nonlinear related pointlike localized field strength. The trick of gauge theory, by which one can extract pointlike localized generators without being required to construct first the string-localized ones, is a resource which does not seem to exist for higher spins, not even if one is willing to cope with unphysical ghosts in intermediate steps. The most interesting interactions are of course the selfinteractions between $(m = 0, s > 1)$. Here one runs into similar problems as with Yang-Mills models (next section). The independence on $e's$ of the local observables leads to nonlinear transformation laws which extend that of free stringlike potentials and the non-existence of linear local observables. Although saying this does not solve any such problem, the lack of an extension of the gauge idea to higher spin makes one at least appreciative of a new view based on localization.

There is one important case which we have left out, namely that of massless Yang-Mills theories interaction with massive matter. This will be commented on in the next section.

There are 2 different categories of delocalization: string-localization with nontrivial pointlike-generated subalgebras as QED. But generically the coupling of string-localized fields leads to a theory with *no local observables*. The models of physical interest are those which contain nontrivial e-independent subfields. For the case at hand the crucial relation is that the change in the string direction can be written as a derivative as in (13). Interactions which are not invariant under law of change of e do not give rise to compact localized observables and physically uninteresting, even if they mathematically exist.

4 Concluding remarks

The guiding model for the presentation of the modular localization alternative to the standard gauge setting has been QED. It would be very interesting to understand the physical consequences of the much stronger string-localization in QCD. The third Wigner class of infinite spin representations is a kinematical illustration of a strong string-localization. Such a string remains invisible to any local or quasilocal registering counter. The reason is that such a counter is limited to register the presence of a local piece of the string but the caused local change would be in contradiction with the holistic nature of an irreducible object. Such objects cannot be created from pointlike or QED charge strings. According to popular argument (without mathematical support) an interaction cannot change the irreducible particle components, if in the interaction density there was no infinite spin component there will be none in the resulting theory.

Whereas for gauge theories the use of string-localized potentials can be considered as a refinement of gauge theory, the use of s=2 string-localized potentials $g_{\mu\nu}(x,e)$ with $d_{sc}=1$ which, without interactions, is linearly related to 4-tensor $R_{\mu\nu\kappa\lambda}$ with the symmetry properties of the Riemann tensor is virgin soil, since in this case there is no gauge setting. Another more speculative idea envisages a of obtaining renormalizable higher massive theories by starting from zero mass and generalizing the screening idea. Since one cannot expect that generic couplings of higher spin massive string-localized potentials lead to pointlike generated subalgebras (not even for s=1), one perhaps needs compensating lower spin fields (in analogy to the screened charged field) in order to select such theories.

Leaving the issue of confinement/invisibility aside, one can certainly study perturbative QCD in the new setting. Since now localization is the central issue, the only compatible method is the Epstein-Glaser iteration. This involves a new problem namely the causal ordering (time-ordering) for strings. For coplanar strings, which are orthogonal to a timelike vector, this can be achieved [13]. Without having done detailed calculation it is not clear whether this lack of covariance is a blessing (the spontaneously broken Lorentz invariance in the charge sectors) or a curse.

For the perturbation theory itself there are two possible strategies. On could either stay close to the spirit of the axial gauge approach;iIn that case one would try to take the composite limit $e_{i_1}....e_{i_k}\rightarrow e_i$ so that the i^{th} charge line has just one string direction e_i which can then be used to control the infrared properties. On the other hand there is the setting of Bogoliubov's S functional which would favor an approach in which the points e on de Sitter space receive the same treatment the $x's$ in Minkowski space. In that case one would integrate over all $e's$ which would seem to totally delocalize the strings. Of course on could try to reintroduce the string-localization by working with string-localized matter fields and integrate, as one does with the x, over all internal e of a graph. The fact tha such a procedure does not speak against it, as long as the S-matrix stays the same. More details and hopefully the resolution of these questions will be contained in a forthcoming collaboration [24].

References

[1] P. Jordan, in: *Talks and Discussions of the Theoretical-Physical Conference in Kharkov* (May 19-25, 1929) Physikalische Zeitschrift **XXX**, (1929) 700-725

[2] B. Schroer, *Pascual Jordan's legacy and the ongoing research in quantum field theory*, arXiv:1010.4431

[3] R. Haag, *Local Quantum Physics*, Springer, second edition, 1996

[4] F. Bloch and A.Nordsiek, Phys. Rev. **52**, (1937) 54

[5] D. Yenni, S. Frautschi and H. Suura, *The infrared divergence phenomena and high energy processes*, Ann. of Phys. **13**, (1961) 370-462

[6] B. Schroer, *Unexplored regions in QFT and the conceptual foundations of the Standard Model*, . arXiv:1010.4431

[7] R. Brunetti, D. Guido and R. Longo, *Modular localization and Wigner particles*, Rev.Math.Phys. **14**, (2002) 759

[8] J. Mund, B. Schroer and J. Yngvason, *String-localized quantum fields and modular localization*, CMP **268** (2006) 621, math-ph/0511042

[9] S. Weinberg, *The Quantum Theory of Fields I*, Cambridge University Press 1995

[10] B. Schroer, *Studies in History and Philosophy of Modern Physics* **41** (2010) 104

[11] B. Schroer, *Modular localization and the d=1+1 formfactor program*, Annals of Physics **295**, (1999) 190

[12] L. Fassarella and B. Schroer, *Wigner particle theory and local quantum physics*, J. Phys. A **35**, (2002) 9123-9164

[13] J. Mund, *String–Localized Quantum Fields, Modular Localization, and Gauge Theories,* New Trends in Mathematical Physics (V. Sidoravicius, ed.), Selected contributions of the XVth Int. Congress on Math.Physics, Springer, Dordrecht (2009), 495

[14] P. Leyland, J. Roberts and D. Testard, Duality For Quantum Free Fields, C.N.R.S. preprint July 1978, unpublished

[15] S. Weinberg, *Feynman rules for any spin*, Phys. Rev. **133**B (1964) 1318

[16] B. L Joffe, V. S. Fadin and L. N. Lipatov, *Quantum Chromodynamics*, Cambridge Monographs on particle Physics, Nuclear Physics and Cosmology, Cambridge University Press 2010

[17] D. Buchholz, Commun. Math. Phys. **52**, (1977), 147

[18] O. Steinmann, *A New Look at the Higgs-Kibble Model*, arXiv:0804.2989

[19] B. Schroer, *Particle physics in the 60s and 70s and the legacy of contributions by J. A. Swieca*, arXiv:0712.0371

[20] J. A. Swieca, Phys. Rev. D **13**, (1976) 312

[21] M. Duetsch and B. Schroer, J.Phys. **A33**, (2000) 4317

[22] M. Dütsch, J. M. Gracia-Bondia, F. Scheck and J. C. Varilly, *Quantum gauge models without classical Higgs mechanism*, arXiv:1001.0932

[23] J. H. Lowenstein and B. Schroer, Phys. Rev. D **6**, (1972) 1553

[24] J- Mund and B. Schroer, work in progress