Off-Shell Supersymmetry

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Supersymmetry does not dictate the way we should quantize the fields in the supermultiplets, and so we have the freedom to quantize the Standard Model (SM) particles and their superpartners differently. We propose an unconventional quantization rule such that a particle can only appear off-shell, while its contributions to quantum corrections behave exactly the same as in the usual quantum field theory. We apply this quantization rule solely to the sparticles in the Minimal Supersymmetric Standard Model (MSSM). Thus sparticles can only appear off-shell. They could be light but would completely escape the direct detection at any experiments such as the LHC. However, our theory still retains the same desirable features of the usual MSSM at the quantum level. For instance, the gauge hierarchy problem is solved and the three MSSM gauge couplings are unified in the usual way. Although direct detection of sparticles is impossible, their existence can be revealed by precise measurements of some observables (such as the running QCD coupling) that may receive quantum corrections from them and have sizable deviations from the SM predictions. Also the current experimental constraints from the indirect sparticle search are still applicable.

\textbf{Introduction.} Supersymmetric (SUSY) extension of the Standard Model (SM) is one of the most promising ways to solve the gauge hierarchy problem in the SM. The Minimal Supersymmetric Standard Model (MSSM) not only offers a solution to the gauge hierarchy problem, but also provides several interesting features for particle physics phenomenology. For example, the three MSSM gauge couplings are beautifully unified at $M_{\text{GUT}} \simeq 2 \times 10^{16}$ GeV, which suggests an interesting paradigm, Grand Unified Theory (GUT). With a conserved $R$-parity, the lightest sparticle (neutralino) is a primary candidate for the dark matter in the Universe. Besides, electroweak symmetry breaking can be triggered radiatively from a large top Yukawa coupling in the presence of SUSY-breaking terms. The SM-like Higgs boson mass is then predicted as a function of soft SUSY-breaking terms. The usual way [6, 7]. These models cannot completely hide the direct signal of sparticles forever. With the substantially improved sensitivity at the LHC Run II, it is conceivable that they will be probed and constrained soon.

Contrary to the SUSY models which can tentatively hide the sparticles from direct detection, we would like to provide an MSSM scenario that can survive even if sparticles are never observed at the LHC. Our guiding principle is to retain simplicity and naturalness as much as possible. Our work is also based on the observation that we do not need on-shell sparticles in order to solve the gauge hierarchy problem and unify the three MSSM gauge couplings.

In this paper, we propose an unconventional quantization rule in which a particle can only appear off-shell. We then apply this rule to quantize the sparticles in the MSSM while the SM particles are quantized in the conventional way. As a result, sparticles can only appear as the off-shell state, which makes their direct detection impossible by any experiments. However, the treatment of quantum corrections in our theory is exactly the same as in the usual quantum field theory (QFT). Thus our MSSM retains the same desirable features in terms of quantum corrections as the usual MSSM. For instance, the gauge hierarchy problem is solved and the three MSSM gauge couplings are unified in the usual manner. The experimental constraints from the indirect sparticle search are still applicable and the same as in the usual MSSM. Therefore, even though sparticles only

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appear off-shell, SUSY should be broken with stringent bounds on flavor and additional CP violations [8].

In the following sections, we first describe our unconventional quantization rule which leads to off-shell particles, and then we resolve some apparent pathological issues associated with it. We impose this unconventional quantization rule on the sparticles in the MSSM, so that all the experimental bounds from the direct sparticle search disappear. Finally, we discuss the phenomenology of off-shell sparticles and possibilities to indirectly detect them.

**Unconventional Quantization.** In the usual QFT, a real scalar field \( \phi \) is quantized in the following way:

\[
\phi(x) = \int \frac{d^3p}{(2\pi)^3/2} \sqrt{\omega_p} \left( a(p) e^{-ipx} + a^\dagger(p) e^{ipx} \right),
\]

where the annihilation operator \( a(p) \) and the creation operator \( a^\dagger(p) \) obey the commutation relation \( [a(p), a^\dagger(p')] = \delta^{(3)}(p - p') \). This ensures that the equal-time canonical quantization rule \( [\phi(x), \phi^\dagger(x')] = i \delta^{(3)}(x - x') \) is satisfied. The vacuum \( |0\rangle \) is chosen such that \( a(p) |0\rangle = 0 \). As we know, if we compute the propagator for \( \phi \), we will obtain the Feynman propagator which is a retarded propagator.

Contrary to the conventional wisdom, we would like to introduce a different quantization rule. We retain the commutation relation \( [a(p), a^\dagger(p')] = \delta^{(3)}(p - p') \), but propose that

\[
a | n \rangle = \text{sign} \left( n - \frac{1}{2} \right) \sqrt{\left| n - \frac{1}{2} \right|} | n - 1 \rangle,
\]

\[
a^\dagger | n \rangle = \sqrt{\left| n + \frac{1}{2} \right|} | n + 1 \rangle,
\]

for any integer \( n \) which characterizes the number of energy units carried by the state \( | n \rangle \). Notice that for simplicity, we have suppressed the dependence on momentum \( p \) in Eq. (2) and Eq. (3). (Apparently, this quantization rule may lead to the disastrous negative energies and probably even worse problems. We will hold on and resolve the apparent pathological issues below Eq. (11). In fact, the resolution may become impossible when we move on.) To elaborate, with the above choice, we have \( a | 0 \rangle = -\sqrt{1/2} | -1 \rangle \) and \( a^\dagger | 0 \rangle = \sqrt{1/2} | 1 \rangle \). This procedure implies that

\[
\langle 0 | a^\dagger(p) a(p') | 0 \rangle = -\frac{1}{2} \delta^{(3)}(p - p'),
\]

\[
\langle 0 | a(p) a^\dagger(p') | 0 \rangle = \frac{1}{2} \delta^{(3)}(p - p').
\]

Thus, the vacuum expectation value of the Hamiltonian for \( \phi \) without normal ordering, \( H_\phi = 1/2 \int d^3p \omega_p \left( a^\dagger(p) a(p) + a(p) a^\dagger(p) \right) \), is

\[
\langle 0 | H_\phi | 0 \rangle = 0.
\]

In other words, the vacuum state \( | 0 \rangle \) is no longer a state with the lowest energy but is simply a state with no particles and thereby zero energy.

Applying the above quantization procedures, the propagator for \( \phi \) turns out to be an average of the usual Feynman propagator with an \( +ive \) prescription and a similar one with an \( -ive \) prescription:

\[
\langle 0 | T\{ \phi(x) \phi(y) \} | 0 \rangle = \frac{1}{2} \left( G_+(x - y) + G_-(x - y) \right),
\]

where \( G_+(x - y) \) and \( G_-(x - y) \) are respectively given by

\[
G_+(x - y) = \int \frac{d^4p}{(2\pi)^4} e^{-ip(x-y)} \frac{i}{p^2 - m^2 + i\epsilon},
\]

\[
G_-(x - y) = \int \frac{d^4p}{(2\pi)^4} e^{-ip(x-y)} \frac{i}{p^2 - m^2 - i\epsilon}.
\]

Using the identity

\[
\frac{1}{p^2 - m^2 \pm i\epsilon} = P \left( \frac{1}{p^2 - m^2} \right) \mp i\pi \delta(p^2 - m^2),
\]

where \( P(\cdots) \) denotes the principal-part, we obtain

\[
\langle 0 | T\{ \phi(x) \phi(y) \} | 0 \rangle
\]

\[
= \int \frac{d^4p}{(2\pi)^4} e^{-ip(x-y)} P \left( \frac{i}{p^2 - m^2} \right).
\]

Therefore, \( \phi \) acquires a principal-value propagator instead of the usual Feynman propagator. Actually, one can verify that this propagator is equivalent to an average of the retarded and the advanced Green’s functions for a scalar field. Thus, it is a half retarded and half advanced propagator. Since there is no longer a support from the delta function \( \delta(p^2 - m^2) \) in the propagator, the \( \phi \) particle cannot be on-shell. In other words, the \( \phi \) particle is “ghost-like” in the sense that it can only appear off-shell. However, the procedure of making the \( \phi \) particle “ghost-like” is completely different from the usual Faddeev-Popov approach [9]. We do not need a ghost Lagrangian. What we need is simply the unconventional quantization rule described above.

Since the \( \phi \) particle can only appear off-shell, the possibility of negative energies does not cause any problem. In fact, one can show that if the state \( | n \rangle \) has a negative energy with \( n < 0 \), it may also acquire a negative norm:

\[
\langle n | n \rangle < 0, \quad \text{if } n < 0 \text{ and } | n | = \text{odd},
\]

which implies the existence of an indefinite metric. But again, since the \( \phi \) particle can only appear off-shell, it does not contribute to the unitarity sum of the scattering amplitudes. Hence, unitarity is preserved despite the possible existence of negative-norm states. This is analogous to the Faddeev-Popov ghosts which have negative norms but can only appear off-shell. The idea of indefinite metric was first invented by Dirac [10] and then
elaborated by Pauli [11] in 1940s in an attempt to remove the divergences and construct a finite theory of quantum electrodynamics. Their attempt turned out to be not very satisfactory. But the canceling effect due to indefinite metric inspired Lee and Wick to construct an alternative finite theory of quantum electrodynamics [12]. The possible issues of Lee-Wick theory were discussed in [13]. For a comprehensive lecture and review on indefinite metric QFT, one can consult [14] and [15] respectively. In any case, due to the success of renormalization, these attempts for taming the divergences were largely forgotten.

Since we have retained \([a(p), a^†(p′)] = δ^{(3)}(p − p′)\) in spite of the unconventional rule introduced in Eq. (2) and Eq. (3), it follows that the equal-time commutator \([ϕ(x), ϕ(y)]|_{x^0 = y^0}\) should vanish as usual. This means that micro-causality is preserved under our framework. Moreover, for spacelike distances with \((x − y)^2 = −r^2 < 0\), there exists a reference frame where \(x^0 − y^0 = 0\). It is well-known that

\[
G_+(x − y) = \frac{m}{4π^2 r} K_1(m r), \quad \text{for } (x − y)^2 < 0, \quad (13)
\]

where \(K_1(z)\) is the modified Bessel function of second kind. Thus outside the light-cone, the usual Feynman propagator predicts that the propagation amplitude is exponentially small but nonzero (as \(r \to \infty\)). In contrast, one can show that

\[
G_−(x − y) = −θ(x^0 − y^0) \int \frac{d^3 p}{(2π)^3} \frac{1}{2ω_p} e^{-ip(x − y)}
\]

\[
−θ(y^0 − x^0) \int \frac{d^3 p}{(2π)^3} \frac{1}{2ω_p} e^{-ip(y − x)}, \quad (14)
\]

and hence we have

\[
G_−(x − y) = −G_+(x − y), \quad \text{for } (x − y)^2 < 0. \quad (15)
\]

Therefore, our propagator, which is half-retarded and half-advanced, is identically zero for spacelike distances. In other words, the propagator obtained from our unconventional quantization scheme predicts that the propagation amplitude is exactly zero outside the light-cone and so it is truly causal.

In fact, our half-retarded and half-advanced propagator may have some resemblance to the absorber theory of radiation proposed by Feynman and Wheeler [16]. They considered a half-retarded and half-advanced electromagnetic field in which electrons radiate symmetrically, both forward and backward in time. As discussed in [17], due to subtle cancellations, the absorber theory is “an apparently acausal theory that is not”. Similarly, under our framework, each of the positive-energy and negative-energy states is propagated symmetrically, both forward and backward in time. The subtle “destructive interference” renders the φ particle virtual (off-shell) and, at the same time, causal.

The generalization of our unconventional quantization scheme to a complex scalar field as well as a vector boson is straightforward. (In the MSSM, we do not have spin-1 sparticles, so the generalization to vector bosons is actually irrelevant.) For a complex scalar given by

\[
Φ(x) = \int \frac{d^3 p}{(2π)^3/2√ω_p} (a(p) e^{-ipx} + b^†(p) e^{ipx}),
\]

we require the usual commutation relations \([a(p), a^†(p′)] = [b(p), b^†(p′)] = δ^{(3)}(p − p′)\) and impose that similar to \(a(p)\) and \(a^†(p)\), \(b(p)\) and \(b^†(p)\) acting on the state \(|n\rangle\) satisfy the relations in Eq. (2) and Eq. (3) respectively. The vacuum expectation value of the Hamiltonian for Φ without normal ordering, \(H_Φ = \frac{1}{2} \int d^3 p ω_p (a(p) a^†(p) + a(p) a^†(p) + b(p) b^†(p) + b(p) b^†(p))\), is \(\langle 0 | H_Φ | 0 \rangle = 0\). Besides, the propagator for Φ, \(\langle 0 | T \{ Φ(x) Φ(y) \} | 0 \rangle\), is of the same form as Eq. (11).

We continue our unconventional quantization scheme to a Dirac fermion:

\[
ψ(x) = \int \frac{d^3 p}{(2π)^3/2√2ω_p} \sum_{s = ±} \left( c(p, s) u(p, s) e^{-ipx} + d^†(p, s) v(p, s) e^{ipx} \right).
\]

Here, the creation and annihilation operators obey the usual anticommutation relations \(\{c(p, s), c^†(p′, s′)\} = \{d(p, s), d^†(p′, s′)\} = δ^{(3)}(p − p′) δ_{ss′}.\) The spinors \(u(p, s)\) and \(v(p, s)\) satisfy the usual orthogonality conditions as well as the completeness relations: \(\sum_s u(p, s) u^†(p, s) = ψ + m\) and \(\sum_s v(p, s) v^†(p, s) = ψ − m\). We introduce the following quantization steps:

\[
c \quad \text{or} \quad d \quad n \quad = \sqrt{|n − 1/2|} \quad |n − 1\rangle, \quad (17)
\]

\[
c^† \quad \text{or} \quad d^† \quad n \quad = \sqrt{|n + 1/2|} \quad |n + 1\rangle, \quad (18)
\]

for simplicity, we have suppressed the dependence on momentum \(p\) and spin index \(s\) in Eq. (17) and Eq. (18). These imply that

\[
\langle 0 | c^†(p, s) c(p′, s′) | 0 \rangle = \langle 0 | c(p, s) c^†(p′, s′) | 0 \rangle = \frac{1}{2} δ^{(3)}(p − p′) δ_{ss′}, \quad (19)
\]

\[
\langle 0 | d^†(p, s) d(p′, s′) | 0 \rangle = \langle 0 | d(p, s) d^†(p′, s′) | 0 \rangle = \frac{1}{2} δ^{(3)}(p − p′) δ_{ss′}. \quad (20)
\]

Notice that due to the anticommutation relations, a difference from the bosonic case is that only \(n = −1, 0, 1\) are allowed. One can verify that the norm for each of these states is positive-definite. The vacuum expectation value of the Hamiltonian for ψ without normal ordering, \(H_ψ = \sum_{s = ±1/2} \int d^3 p ω_p (c^†(p, s) c(p, s) − d(p, s) d^†(p, s)),\) is

\[
\langle 0 | H_ψ | 0 \rangle = 0. \quad (21)
\]

Similar to the scalars, this means that the vacuum state \(|0\rangle\) is no longer a state with the lowest energy but is simply a state with no particles and thereby zero energy.
The propagator for $\psi$ is also half-retarded and half-advanced, where (with spinor indices suppressed) the corresponding propagator $D_{-}(x - y)$ with an $-ie$ prescription is given by

$$D_{-}(x - y) = \int \frac{d^{4}p}{(2\pi)^{4}} e^{-ip(x-y)} \frac{i(\not{p} + m)}{p^{2} - m^{2} - i\epsilon} \quad (22)$$

$$= \theta(x^{0} - y^{0}) \int \frac{d^{3}p}{(2\pi)^{3}} \frac{\not{p} - m}{2\omega_{p}} e^{-i\omega_{p}(y-x)} - \theta(y^{0} - x^{0}) \int \frac{d^{3}p}{(2\pi)^{3}} \frac{\not{p} + m}{2\omega_{p}} e^{-i\omega_{p}(x-y)}. \quad (23)$$

The exact form of the propagator for $\psi$ is

$$\langle 0 | T \{ \psi_{\alpha}(x) \bar{\psi}_{\beta}(y) \} | 0 \rangle \equiv S(x - y)_{\alpha\beta}$$

$$= \int \frac{d^{4}p}{(2\pi)^{4}} e^{-ip(x-y)} (\not{p} + m)_{\alpha\beta} P \left( \frac{i}{p^{2} - m^{2}} \right), \quad (24)$$

and so the $\psi$ particle can only appear off-shell. Following the similar arguments for the scalars, the possibility of negative energies is not a problem and micro-causality is preserved. Besides, one can verify, using Eq. (23), that this propagator is exactly zero outside the light-cone.

For a Majorana fermion with $\psi = \psi^{c} = C \bar{\psi}^{T}$ where $C$ is the charge conjugation matrix, the quantization steps are similar except that we set $c(p, s) = d(p, s)$. There are two additional propagators:

$$\langle 0 | T \{ \lambda_{a}(x) \psi_{\beta}(y) \} | 0 \rangle = \left[ C^{-1} S(x - y) \right]_{a\beta} \quad (25)$$

$$\langle 0 | T \{ \bar{\lambda}_{a}(x) \bar{\psi}_{\beta}(y) \} | 0 \rangle = \left[ -S(x - y) C \right]_{a\beta}. \quad (26)$$

**SUSY with Unconventional Quantization.** For a given supermultiplet, supersymmetry exhibited at the Lagrangian does not constrain how we quantize the fields. Providing the usual MSSM Lagrangian, we propose that the SM particles are quantized in the conventional way, while their superpartners are quantized according to the unconventional rule described above. Therefore, sparticles can only appear off-shell in our theory, which evades the direct detection at any experiments. However, since quantum corrections are calculated in the same way as in the usual QFT, the gauge hierarchy problem is solved in the usual way.

Although all the experimental constraints from the direct detection disappear, the results from the indirect sparticle search are still applicable to our MSSM scenario to constrain the sparticle mass spectrum. Similar to the usual MSSM, this requires SUSY to be broken with stringent bounds on flavor and additional CP violations. With the soft SUSY-breaking terms, our MSSM scenario with unconventional quantization leads to the same attractive phenomenological consequences as the usual MSSM does [18]. For example, the three MSSM gauge couplings are successfully unified at $M_{\text{GUT}} \approx 2 \times 10^{16}$ GeV, the electroweak symmetry is radiatively broken, and the SM-like Higgs boson mass is obtained through stop loop corrections.

Due to the unconventional quantization method, our SUSY theory has an unusual property. As a consequence of applying the unconventional quantization rule to sparticles, we obtain $\langle 0 | H_{\text{sparticles}} | 0 \rangle = 0$. This implies that $\langle 0 | H_{\text{SM}} + H_{\text{sparticles}} | 0 \rangle \neq 0$ even if SUSY is manifest. Thus, it appears that SUSY is fundamentally broken by the unconventional quantization method. However, the structure of SUSY theory for quantum corrections such as the cancellations of quadratic and quartic divergences is still retained, and so the SUSY breaking in vacuum energy has no practical effect on particle physics phenomenology. We can simply subtract out the constant vacuum energy as in the usual QFT.

**Phenomenology of Off-Shell SUSY.** As mentioned above, although it is impossible to directly detect sparticles which only appear off-shell, the existence of sparticles can be indirectly identified through their contributions to quantum corrections for some observables. For example, the fine structure constant at the Z-pole $(m_{Z})$ is very precisely measured as $\alpha_{\text{em}}(m_{Z})^{-1} = 127.918 \pm 0.019$ [19], which is consistent with the SM prediction for the evolution of the fine structure constant from low energy to the Z-pole. If charged sparticles such as squarks, sleptons and charginos are involved in the evolution, the resultant fine structure constant at the Z-pole will be altered from the SM prediction. This sets the lower bound on the charged sparticle masses as $m \gtrsim m_{Z}$.

The discussion about the fine structure constant may give us an idea about how to identify the existence of light colored (off-shell) sparticles such as the gluino and squarks by the LHC, even though the experimental data would show no indication of sparticle productions. At energies higher than their masses, the colored sparticles are involved in the running QCD coupling and will deflect the running from the trajectory predicted by the SM. Employing the 1-loop renormalization group equation for the QCD coupling, we define a deviation of the running QCD coupling at an energy scale $\mu$ as

$$\Delta(\mu) \equiv \frac{\alpha_{s}^{\text{MSSM}}(\mu)}{\alpha_{s}^{\text{SM}}(\mu)} - 1 \quad (27)$$

$$\approx \frac{\alpha_{s}^{\text{SM}}(m_{\text{Higgs}}) - \alpha_{s}^{\text{SM}}(m_{\ast})}{2\pi} \ln \left( \frac{\mu}{m_{\ast}} \right), \quad (28)$$

where $\alpha_{s}^{\text{SM}}$ is the running SM QCD coupling, $\alpha_{s}^{\text{MSSM}}$ is the running QCD coupling with the contributions from colored sparticles with a degenerate mass $\tilde{m}$, and $b_{\text{SM}} = -7 \cdot (b_{\text{MSSM}}) = -7 \cdot (b_{\text{MSSM}})$ is the QCD beta function coefficient in the SM (MSSM). Using $\alpha_{s}(M_{t}) = 0.0928$ with a top quark pole mass $M_{t} = 173.34$ GeV [20], we find $\Delta(1 \text{ TeV}) \approx 3.6\%$ for $\tilde{m} = 500$ GeV and $b_{\text{MSSM}} = -3$ due to the contributions from the degenerate gluino and three generations of squarks. This deviation is slightly smaller than the error of the current measurement of the QCD coupling constant in the TeV range [21]. In fact, we obtain $\Delta(\mu \gtrsim 3 \text{ TeV}) \gtrsim 9\%$ for the same parameters. Therefore, a more precise measurement of the QCD coupling at sufficiently high energy and luminosity may reveal the off-shell colored sparticles.
Conclusions and Discussions. In this paper, we propose a novel MSSM scenario where sparticles are quantized in an unconventional way to only appear off-shell. However, calculations of quantum corrections in this theory are exactly the same as in the usual QFT. As a consequence, most of the phenomenologically attractive properties of the usual MSSM, such as the solution to the gauge hierarchy problem and the successful gauge coupling unification, remain the same. Although direct detection of sparticles is impossible, their existence can be revealed through precise measurements of observables to which off-shell sparticles give sizable quantum corrections.

In principle, one could apply our unconventional quantization rule to any other theories so as to evade the corresponding bounds from direct detection. Nevertheless, sparticles are particularly well-suited for the unconventional quantization. The reason is that the most important merit of sparticles is due to their off-shell quantum contributions. For instance, we do not need on-shell sparticles in order to solve the gauge hierarchy problem and unify the three MSSM gauge couplings.

Of course, the lightest sparticle (neutralino) would no longer be a viable dark matter candidate if it can only appear off-shell. However, this should not be considered as a serious deficiency of our MSSM scenario. Providing a viable dark matter candidate is a just bonus of the usual MSSM, and it is easy enough to construct other models that give a promising dark matter candidate. Supersymmetry is most crucial for solving the gauge hierarchy problem and unification of the three MSSM gauge couplings.

Also, it is in principle not necessary for all sparticles to be quantized in the unconventional way. For example, as in the usual MSSM, we may apply the conventional quantization rule to the lightest sparticle (neutralino) such that it remains a viable dark matter candidate. If some sparticles are discovered at the LHC Run II, it means that these sparticles are quantized in the conventional way, but it is conceivable that other sparticles are unconventionally quantized.

Our encouraging message is that even if none of sparticles is observed at the LHC Run II, there is still a hope that the sparticles are light but can only appear off-shell. In that case, their existence may have to be indirectly identified through a precise measurement of the running QCD coupling.

In fact, our work is more general than supersymmetry. Our idea has provided a new vision that some new physics may only appear off-shell. An intriguing collider signature of this kind of new physics is that we may see sizable deviations from the SM predictions at precision measurements despite the absence of new particles at direct detection.

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