

Strings, world-sheet covariant quantization and Bohmian mechanics

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Abstract. The covariant canonical method of quantization based on the De Donder–Weyl covariant canonical formalism is used to formulate a world-sheet covariant quantization of bosonic strings. To provide the consistency with the standard non-covariant canonical quantization, it is necessary to adopt a Bohmian deterministic hidden-variable equation of motion. In this way, string theory suggests a solution to the problem of measurement in quantum mechanics.

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1 Introduction

String theory [1–3] is a theory with an ambition to be the theory of everything. However, there is one fundamental problem on which, so far, string theory has had nothing new to say. This is the problem of the interpretation of quantum mechanics (QM), or, in more physical terms, the problem of measurement in QM. From this point of view, it is widely believed that it does not matter whether one quantizes a particle, a field, or a string; the formalism of quantization is, essentially, always the same, so the interpretation adopted, say, for particles, should work equally well (or badly) for fields or strings. Although the recent progress in understanding the phenomenon of decoherence shed much light on the problem of measurement in QM, this problem is still considered unsolved [4]. In this paper, however, we argue that string theory offers a new insight into the problem of interpretation/measurement in QM, an insight that cannot be inferred from the quantization of a particle. Since strings, unlike particles, are extended objects, the requirement of world-sheet covariance leads to a non-trivial relation between the σ^0 -dependence and the σ^1 -dependence of the string coordinates $X^\alpha(\sigma^0, \sigma^1)$. In order to preserve the world-sheet covariance at the quantum level, we argue that the classical covariant De Donder–Weyl canonical formalism (see e.g. [5, 6] and references therein) might be a good starting point. The appropriate quantum formalism is developed in [7] for fields. In particular, the formalism attributes a new status to the Bohmian deterministic hidden-variable interpretation of QM [8–14], because, in [7], the Bohmian equations of motion for fields are derived from the requirement of spacetime covariance. By replacing the requirement of spacetime covariance for

fields with that of world-sheet covariance for strings, in this paper we observe that a completely analogous argument leads to the Bohmian formulation of quantum strings. (In the Bohmian interpretation, the quantum string coordinates $X^\alpha(\sigma^0, \sigma^1)$ evolve in a deterministic manner even when they are not measured.) Thus, in contrast to particle physics where the Bohmian deterministic interpretation is just one of many interpretations of QM, we argue that in string theory the Bohmian interpretation emerges naturally from the requirement of world-sheet covariance.

To further motivate the analysis presented in the subsequent sections, it is worthwhile to explain the conceptual difference between the physical meaning of the results obtained in [7] and that of the present paper. For that purpose, we need to recapitulate the concepts of particles, fields and strings in a somewhat wider context. In non-string theories, quantum fields are often viewed in two different ways. The prevailing point of view among “hard-core” field theorists is that fields are the only fundamental objects, while particles are merely emergent objects that sometimes even cannot be well defined (see e.g. [15–18]). On the other hand, particle-physics phenomenologists are more willing to view pointlike particles as the fundamental objects, while the field is often viewed among them merely as a calculational tool convenient for treating interactions in which the number of particles changes. Indeed, there exists an alternative string-inspired particle-scattering formalism that completely avoids any referring to fields [19]. In string theory, the situation is similar, but with a difference consisting of the fact that most of the work in string theory is done without referring to string-field theory. Moreover, there are indications that string-field theory might not be the correct way to treat string interactions [20]. Thus, from the string-theory perspective, particles might be more fundamental objects than fields.

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In the context of the Bohmian hidden-variable interpretation of QM, the field-or-particle dilemma is even sharper than in the conventional interpretation. Should the Bohmian interpretation be applied to particles, to fields, or to both? Since the conventional probabilistic interpretation cannot be applied to the relativistic Klein–Gordon equation, the Bohmian deterministic interpretation of relativistic quantum particles might be a natural choice with interesting measurable predictions [21, 22]. However, a derivation of the Bohmian interpretation from the requirement of relativistic covariance based on the De Donder–Weyl formalism [7] works for fields, but not for particles. Thus, if particles are more fundamental objects than fields, then the results of [7] might be physically irrelevant and we still cannot derive the Bohmian interpretation. However, now comes string theory to save the situation. If particles are more fundamental than fields, but if they are not really pointlike, but extended objects as in string theory, then the results of [7] can be applied. In this case, the Bohmian interpretation of strings can be derived from the requirement of world-sheet covariance, while the resulting string theory in a pointlike-particle limit reduces to the Bohmian interpretation of relativistic quantum particles.

The classical De Donder–Weyl formalism for bosonic strings is presented in Sect. 2, while the corresponding quantum theory of bosonic strings is formulated in Sect. 3. The case of supersymmetric strings is still beyond our current technical achievements.

2 Classical De Donder–Weyl formalism for bosonic strings

In order to have a notation similar to that in [7], let the letters $\alpha, \beta = 0, 1, \dots, D-1$ denote the target spacetime indices, and let the letters $\mu, \nu = 0, 1$ denote the world-sheet indices. The signature of the spacetime metric is chosen to be $(+, -, \dots, -)$. Similarly, on a flat world-sheet we have $\eta^{00} = -\eta^{11} = 1$. We also use the notation $\sigma \equiv (\sigma^0, \sigma^1)$. With this notation, the action of a bosonic string is

$$A = \int d^2\sigma \mathcal{L}, \quad (1)$$

where

$$\mathcal{L} = -\frac{1}{2} |h|^{1/2} h^{\mu\nu} \eta_{\alpha\beta} (\partial_\mu X^\alpha) (\partial_\nu X^\beta) \quad (2)$$

is the Lagrangian density. Here $\eta_{\alpha\beta}$ is a flat Minkowski metric in D dimensions, $h^{\mu\nu}(\sigma)$ is an arbitrary metric on the string world-sheet, and h is the determinant of $h_{\mu\nu}$. The spacetime and world-sheet indices are raised (lowered) by $\eta^{\alpha\beta}$ ($\eta_{\alpha\beta}$) and $h^{\mu\nu}$ ($h_{\mu\nu}$), respectively. By requiring that the variation of (1) with respect to $h^{\mu\nu}$ should vanish, one obtains that $h_{\mu\nu}$ must be proportional to the induced metric on the world-sheet [3], i.e.

$$h_{\mu\nu}(\sigma) = f(\sigma) (\partial_\mu X^\alpha) (\partial_\nu X_\alpha), \quad (3)$$

where $f(\sigma)$ is an arbitrary positive-valued function.

The canonical momentum world-sheet vector density is defined as

$$\mathcal{P}_\alpha^\mu = \frac{\partial \mathcal{L}}{\partial (\partial_\mu X^\alpha)} = -|h|^{1/2} \partial^\mu X_\alpha. \quad (4)$$

The covariant De Donder–Weyl Hamiltonian density is given by the Legendre transform

$$\begin{aligned} \mathcal{H} &= \mathcal{P}_\alpha^\mu \partial_\mu X^\alpha - \mathcal{L} \\ &= -\frac{1}{2} \frac{h_{\mu\nu}}{|h|^{1/2}} \eta^{\alpha\beta} \mathcal{P}_\alpha^\mu \mathcal{P}_\beta^\nu. \end{aligned} \quad (5)$$

When (4) is satisfied, then $\mathcal{H} = \mathcal{L}$. The covariant Hamilton equations of motion are

$$\partial_\mu X^\alpha = \frac{\partial \mathcal{H}}{\partial \mathcal{P}_\alpha^\mu}, \quad \partial_\mu \mathcal{P}_\alpha^\mu = -\frac{\partial \mathcal{H}}{\partial X^\alpha}. \quad (6)$$

Using (5), we see that the first equation in (6) is equivalent to (4). Since \mathcal{H} in (5) does not depend on X^α , the second equation in (6) leads to the covariant string-wave equation

$$\partial_\mu \left(|h|^{1/2} \partial^\mu X_\alpha \right) = 0. \quad (7)$$

Thus, the classical De Donder–Weyl covariant canonical formalism is equivalent to the classical Lagrangian formalism which also leads to the covariant equation of motion (6). Similarly, it is also equivalent to the ordinary non-covariant Hamilton formalism, in which the Hamiltonian is defined such that only $\mu = 0$ contributes in the first line of (5).

The next step is to introduce the covariant De Donder–Weyl Hamilton–Jacobi formalism. We introduce a vector-density function $\mathcal{S}^\mu(X(\sigma), \sigma)$ that satisfies the De Donder–Weyl Hamilton–Jacobi equation

$$\mathcal{H} + \partial_\mu \mathcal{S}^\mu = 0. \quad (8)$$

Here \mathcal{H} is given by (5) with the replacement

$$\mathcal{P}_\alpha^\mu \rightarrow \frac{\partial \mathcal{S}^\mu}{\partial X^\alpha}. \quad (9)$$

The partial derivative ∂_μ acts only on the second argument of $\mathcal{S}^\mu(X(\sigma), \sigma)$. The corresponding total derivative is given by

$$d_\mu = \partial_\mu + (\partial_\mu X^\alpha) \frac{\partial}{\partial X^\alpha}. \quad (10)$$

For a given solution $\mathcal{S}^\mu(X, \sigma)$ of the De Donder–Weyl Hamilton–Jacobi equation, the σ -dependence of $X^\alpha(\sigma)$ is determined by the equation of motion

$$-|h|^{1/2} \partial^\mu X_\alpha = \frac{\partial \mathcal{S}^\mu}{\partial X^\alpha}. \quad (11)$$

The classical De Donder–Weyl Hamilton–Jacobi formalism above has a manifest world-sheet covariance. We would like to construct an analogous quantum formalism with a manifest world-sheet covariance at the quantum

level. It is already known how to construct the quantum formalism that corresponds to the ordinary non-covariant Hamilton–Jacobi formalism: by using quantum mechanics represented by the Schrödinger equation. Thus, the first step towards quantization based on the covariant De Donder–Weyl Hamilton–Jacobi formalism is to explain how the ordinary non-covariant Hamilton–Jacobi formalism can be obtained from the covariant one. Choosing $h_{\mu\nu} = \eta_{\mu\nu}$, (8) can be written in an explicit form:

$$-\frac{1}{2} \frac{\partial \mathcal{S}^0}{\partial X^\alpha} \frac{\partial \mathcal{S}^0}{\partial X_\alpha} + \frac{1}{2} \frac{\partial \mathcal{S}^1}{\partial X^\alpha} \frac{\partial \mathcal{S}^1}{\partial X_\alpha} + \partial_0 \mathcal{S}^0 + \partial_1 \mathcal{S}^1 = 0. \quad (12)$$

Using (10) and (11), the last term can be written as

$$\partial_1 \mathcal{S}^1 = d_1 \mathcal{S}^1 - (\partial_1 X^\alpha)(\partial_1 X_\alpha). \quad (13)$$

Similarly, the second term in (12) can be written as $(1/2)(\partial_1 X^\alpha)(\partial_1 X_\alpha)$. Now we introduce the quantity

$$S = \int d\sigma^1 \mathcal{S}^0, \quad (14)$$

so that

$$\frac{\partial \mathcal{S}^0(X(\sigma), \sigma)}{\partial X^\alpha(\sigma)} = \frac{\delta S([X(\sigma^0), \sigma^1], \sigma^0)}{\delta X^\alpha(\sigma^1; \sigma^0)}, \quad (15)$$

where

$$\frac{\delta}{\delta X^\alpha(\sigma^1; \sigma^0)} \equiv \frac{\delta}{\delta X^\alpha(\sigma^1)} \Big|_{X(\sigma^1)=X(\sigma)} \quad (16)$$

is the functional derivative. Thus, by integrating (12) over $d\sigma^1$, we obtain the ordinary non-covariant Hamilton–Jacobi equation

$$H + \partial_0 S = 0, \quad (17)$$

where

$$H = - \int d\sigma^1 \left[\frac{1}{2} \frac{\delta S}{\delta X^\alpha(\sigma^1; \sigma^0)} \frac{\delta S}{\delta X_\alpha(\sigma^1; \sigma^0)} + \frac{1}{2} (\partial_1 X^\alpha)(\partial_1 X_\alpha) \right] \quad (18)$$

is written for the σ^0 -dependent string coordinate $X^\alpha(\sigma^0, \sigma^1)$. The integral of the total derivative $\int d\sigma^1 d_1 \mathcal{S}^1$ is ignored because it is a constant without any physical significance. The σ^0 -evolution of $X^\alpha(\sigma^0, \sigma^1)$ is given by

$$-\partial^0 X_\alpha(\sigma^0, \sigma^1) = \frac{\delta S}{\delta X^\alpha(\sigma^1; \sigma^0)}, \quad (19)$$

which is a consequence of the $\mu = 0$ component of (11). The covariant constraint (3) implies the non-covariant Hamiltonian constraint $H = 0$ [3].

To anticipate the implications to the quantum case, here it is crucial to observe the following. First, to derive (17) from (8), it was necessary to use the $\mu = 1$ component of (11). Second, if the world-sheet covariance is required, then the validity of the $\mu = 1$ component of (11)

also implies the validity of the $\mu = 0$ component of (11). Third, the validity of the $\mu = 0$ component of (11) implies the classical determinism encoded in (19). Thus, the determinism in classical string theory can be derived from the world-sheet covariance and the requirement that the covariant Hamilton–Jacobi equation (8) and the non-covariant one (17) should be both valid. As we shall see in the next section, a similar argument leads to a derivation of the Bohmian deterministic hidden-variable formulation of quantum strings.

3 Quantization and Bohmian mechanics

How to quantize strings such that the world-sheet covariance is manifest? The standard method is the path-integral quantization based on calculating the generating functional $Z = \int [dX][dh] \exp(iA/\hbar)$. (To avoid an anomaly, one must fix $D = 26$ [23].) This method is useful for calculating Green functions and scattering amplitudes.

Although this is usually sufficient for calculating quantities that are measured in practice, there are also quantities that can be measured in principle but cannot be calculated in a covariant way from Z . In particular, the generating functional Z does not describe a quantum state at a given time. Thus, certain physical information is not described by the path-integral quantization. In order to obtain such information, one can try to use the σ^0 -dependent quantum states $\Psi([X(\sigma^1)], \sigma^0)$ that satisfy the functional Schrödinger equation

$$\hat{H}\Psi = i\hbar\partial_0\Psi, \quad (20)$$

where

$$\hat{H} = - \int d\sigma^1 \left[\frac{-\hbar^2}{2} \frac{\delta}{\delta X^\alpha(\sigma^1)} \frac{\delta}{\delta X_\alpha(\sigma^1)} + \frac{1}{2} (\partial_1 X^\alpha)(\partial_1 X_\alpha) \right]. \quad (21)$$

However, not all states satisfying (20) are physical. In particular, physical states satisfy the Hamiltonian constraint

$$(\hat{H} + a)\Psi = 0, \quad (22)$$

where a originates from a constant that can be added to the action (1) without changing the classical properties of strings. A more common view of this constant is in terms of an operator-ordering constant that can be fixed uniquely [1–3]. The discussion of the value of a as well as the discussion of other requirements on physical states related to the requirement of target spacetime covariance are beyond the scope of the present paper. We only note that (22) and (20) imply that all physical states have the same trivial dependence on σ^0 , i.e. that $\Psi([X(\sigma^1)], \sigma^0) = \Psi[X(\sigma^1)]e^{ia\sigma^0/\hbar}$.

To write (20) and (21), one has to fix a special world-sheet coordinate σ^0 . However, any such choice breaks the world-sheet covariance. To solve this problem, we want

to find a covariant substitute for the Schrödinger equation (20). The similarity of the Schrödinger equation (20) to the non-covariant Hamilton–Jacobi equation (17) suggests that a covariant substitute for (20) might be an equation similar to the covariant De Donder–Weyl Hamilton–Jacobi equation (8). Indeed, the general method of quantization based on the De Donder–Weyl Hamilton–Jacobi equation is developed in [7]. (For a different method, with problems discussed in [7], see also [24, 25].) Here we apply the general results of [7] to the case of bosonic strings.

The first step is to write

$$\Psi = Re^{iS/\hbar}, \quad (23)$$

where R and S are real functionals. One finds that the complex equation (20) is equivalent to a set of two real equations

$$-\int d\sigma^1 \left[\frac{1}{2} \frac{\delta S}{\delta X^\alpha(\sigma^1)} \frac{\delta S}{\delta X_\alpha(\sigma^1)} + \frac{1}{2} (\partial_1 X^\alpha)(\partial_1 X_\alpha) - \mathcal{Q} \right] + \partial_0 S = 0, \quad (24)$$

$$-\int d\sigma^1 \left[\frac{1}{2} \frac{\delta R}{\delta X^\alpha(\sigma^1)} \frac{\delta S}{\delta X_\alpha(\sigma^1)} - \mathcal{J} \right] + \partial_0 R = 0, \quad (25)$$

where

$$\mathcal{Q} = \frac{\hbar^2}{2R} \frac{\delta^2 R}{\delta X^\alpha(\sigma^1) \delta X_\alpha(\sigma^1)}, \quad \mathcal{J} = -\frac{R}{2} \frac{\delta^2 S}{\delta X^\alpha(\sigma^1) \delta X_\alpha(\sigma^1)}. \quad (26)$$

We see that (24) is very similar to (17) with (18), differing from it only in containing the additional quantum \mathcal{Q} -term.

Now, following [7], we replace the classical De Donder–Weyl Hamilton–Jacobi equation (8) with the quantum one,

$$-\frac{1}{2} \frac{h_{\mu\nu}}{|h|^{1/2}} \eta^{\alpha\beta} \frac{d\mathcal{S}^\mu}{dX^\alpha} \frac{d\mathcal{S}^\nu}{dX^\beta} + \mathcal{Q} + \partial_\mu \mathcal{S}^\mu = 0. \quad (27)$$

Here $\mathcal{S}^\mu([X], \sigma)$ is a functional of $X(\sigma)$ and a function of σ , which incorporates quantum nonlocalities in a covariant manner. The derivative d/dX^α is a generalization of the derivative $\partial/\partial X^\alpha$, such that the action of the derivative on nonlocal functionals is well defined [7]. The quantum potential \mathcal{Q} is defined as in (26), but with the replacement $\delta/\delta X^\alpha \rightarrow \delta/\delta_C X^\alpha$. The derivative $\delta/\delta_C X^\alpha$ is a covariant version of the derivative (16). The label C denotes a curve on the world-sheet that generalizes the curve $\sigma^0 = \text{constant}$ in (16). The foliation of the world-sheet into curves C is induced by the dynamical vector density $\mathcal{R}^\mu([X], \sigma)$; the curves are defined by requiring that \mathcal{R}^μ should be orthogonal to the curves at each point σ . The vector density \mathcal{R}^μ satisfies the dynamical equation of motion

$$-\frac{1}{2} \frac{h_{\mu\nu}}{|h|^{1/2}} \eta^{\alpha\beta} \frac{d\mathcal{R}^\mu}{dX^\alpha} \frac{d\mathcal{S}^\nu}{dX^\beta} + \mathcal{J} + \partial_\mu \mathcal{R}^\mu = 0, \quad (28)$$

where \mathcal{J} is defined as in (26) with $\delta/\delta X^\alpha \rightarrow \delta/\delta_C X^\alpha$. The functionals R and S are defined in a covariant way by

$$R = \int_C d\Sigma_\mu R^\mu, \quad S = \int_C d\Sigma_\mu S^\mu, \quad (29)$$

where $R^\mu = |h|^{-1/2} \mathcal{R}^\mu$ and $S^\mu = |h|^{-1/2} \mathcal{S}^\mu$ transform as vectors. In the measure $d\Sigma_\mu = dl n_\mu$, dl is an element of the invariant length of C , while n_μ is a unit vector orthogonal to C . Note that the second equation in (29) is a covariant version of (14). The functionals R and S in (29) define the wave functional Ψ as in (23).

From the covariant formalism above, the non-covariant Schrödinger equation (20) can be derived as a special case. Assume that $\mathcal{R}^\mu = (\mathcal{R}^0, 0)$, that \mathcal{S}^1 is a local functional, and that \mathcal{R}^0 and \mathcal{S}^0 are functionals local in the coordinate σ^0 (see [7] for the precise definitions of these notions of locality!). Then, similarly as in the classical case, by choosing $h_{\mu\nu} = \eta_{\mu\nu}$ and integrating (27) and (28) over $d\sigma^1$, one recovers (24) and (25), which, in turn, are equivalent to the Schrödinger equation (20). (The constant originating from the integral $\int d\sigma^1 d_1 \mathcal{S}^1$ can be absorbed into the constant a .) Just as in the classical case, to obtain (24) from (27), it is necessary to assume that a quantum analog of the $\mu = 1$ component of (11) is valid. The covariance then implies that the $\mu = 0$ component is also valid, so we have a covariant quantum relation

$$-|h|^{1/2} \partial^\mu X_\alpha = \frac{d\mathcal{S}^\mu}{dX^\alpha}. \quad (30)$$

The $\mu = 0$ component of (30) implies that the non-covariant Schrödinger equation (20) should be supplemented with (19). In the quantum context, (19) is nothing but the Bohmian deterministic equation of motion for the σ^0 -dependent hidden variable $X^\alpha(\sigma^0, \sigma^1)$. Indeed, by analogy with the Bohmian interpretation of particles and fields [8–14], (19) could have been postulated immediately after writing (20), as an equation that provides a consistent Bohmian deterministic hidden-variable interpretation of quantum strings. In this interpretation, the wave function is a physical object which does not “collapse” during measurements. The nonlocality encoded in the wave function reflects in a nonlocal quantum potential \mathcal{Q} , which provides nonlocalities needed for a hidden-variable theory to be consistent with the Bell theorem. In the deterministic Bohmian interpretation, all quantum uncertainties are an artefact of the ignorance of the actual initial conditions $X^\alpha(\sigma^1)$ at some initial σ^0 . For more details on this interpretation, we refer the reader to the seminal work [8] and reviews [9–12]. Here, however, the crucial equation of the Bohmian interpretation, (19), is not postulated for interpretational purposes, but derived from the requirement of world-sheet covariance! To be more precise, we stress that the covariant quantum (27) and (28) by themselves do not imply the determinism covariantly encoded in (30). Instead, the need for the determinism encoded in (30) emerges from the requirement that these covariant equations should be compatible with standard non-covariant canonical quantum equations.

Equations (27) and (30) imply a quantum version of (7), namely

$$\partial_\mu (|h|^{1/2} \partial^\mu X_\alpha) - \frac{d\mathcal{Q}}{dX^\alpha} = 0. \quad (31)$$

The covariant quantum constraint takes the same form as the classical one (3). (Note that (3) would be meaningless

in the conventional interpretation of the Schrödinger picture, that does not attribute a definite dependence on σ^0 to X^α .) The non-covariant quantum constraint (22) can be derived from the covariant one in a similar way as (24) and (25) are derived from (27) and (28), provided, in addition, that a constant is added to the action and that \mathcal{R}^0 does not explicitly depend on σ^0 .

We also note that the constraint (22) in the pointlike-particle limit reduces to the massless Klein–Gordon equation, provided that $a = 0$ in the pointlike-particle limit. (A heuristic way to obtain $a = 0$ in the pointlike-particle limit is to recall [1–3] that, in bosonic string theory, a turns out to be proportional to $\sum_{n=0}^{\infty} n$, which leads to a finite value after the analytic continuation of the zeta function. Since n represents the wave-mode number of a string, only $n = 0$ contributes in the pointlike-particle limit, which leads to $a = 0$.) The Bohmian equation of motion (19) leads to

$$\frac{dX_\alpha}{ds} = -\frac{\partial S}{\partial X^\alpha} \quad (32)$$

in the pointlike-particle limit, where $s \equiv \sigma^0$. This can be viewed as a stringy derivation of the relativistic-covariant Bohmian interpretation of the massless Klein–Gordon equation, which, in turn, leads to interesting measurable predictions [22].

Of course, string theory can also describe particles with arbitrary spin, not by considering the pointlike limit, but by considering string states Ψ with fixed quantum numbers that determine the spin (see e.g. [3]). By integrating (19) over $d\sigma^1$, one obtains the Bohmian equation of motion (32) for a particle with an arbitrary integer spin. If bosonic strings are replaced with superstrings, then the Bohmian interpretation of half-integer spin particles can also be included in the same way. However, we do not know yet how to formulate the quantum de Donder–Weyl formalism for supersymmetric strings, so we have not yet been able to derive the Bohmian equation of motion for all spins from the requirement of world-sheet covariance. Our current technical achievements allow one only to derive the Bohmian equation of motion for particles with integer spin.

To summarize, in this paper we have used a new manifestly covariant canonical method of quantization developed in [7] to quantize bosonic strings in a way that provides manifest world-sheet covariance. This new method of quantization, based on the classical De Donder–Weyl covariant canonical formalism, is more general than the standard non-covariant canonical quantization in the Schrödinger picture. The covariant method of quantization contains the non-covariant one as a special case (see also [7]

for a discussion of this point). From the requirement that the covariant method of quantization should lead to the standard non-covariant quantization without violating covariance, it turns out that the quantization method should be supplemented with an equation that corresponds to the Bohmian deterministic hidden-variable formulation of QM. Thus, string theory together with the new quantization method proposed in [7] offers a new insight into the problem of interpretation and measurement in QM (an insight that cannot be inferred from the quantization of a particle) by deriving the Bohmian interpretation from the requirement of world-sheet covariance.

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